

# The Quark-Gluon Plasma - A Short Introduction

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## 1. States of Matter

The quark-gluon plasma is a state of strongly interacting matter, in which the quarks and gluons, which make up hadrons, are not longer confined to color-neutral entities of hadronic size. What does that mean?

Matter, in statistical mechanics, is a system of many constituents in local thermal equilibrium – i.e., a system whose average properties are specified by a few global observables (temperature, energy density, net “charge”). For different values of these observables, the system may exhibit fundamentally different average properties, and so there exist different states of matter, with “phase” transitions occurring when the system changes from one state to the other. The classical pattern, already proposed in both Greek and Hindu natural philosophy, has the form shown in Fig. 1. To four fundamental states: solid, liquid, gas and plasma, the vacuum is added as fifth element (“quintessence”), providing the space in which matter exists. What are the states of matter in the sub-atomic world of strong interactions?

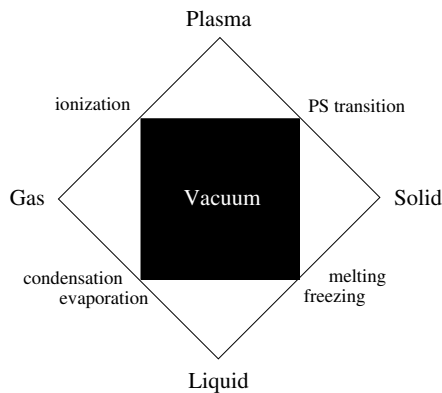


Figure 1: Classical states of matter and transitions between them

To get a first idea, let us begin with a very simple picture. If nucleons, with their given spatial extension, were both elementary and incompressible, then a state of close packing would constitute the high density limit of matter. If, on the other hand, nucleons are really composite – bound states of point-like quarks –, then with increasing density they will start to overlap, until eventually we reach a state in which each quark finds within its immediate vicinity a considerable number of other quarks. It has no way to identify which of these had been its partners in a specific

nucleon at some previous state of lower density. Beyond a certain point, the concept of a hadron thus loses its meaning, and we are quite naturally led from nuclear matter to a system whose basic constituents are unbound quarks.

More specifically, in confined matter the constituents are color-neutral quark-antiquark or three-quark states of hadronic size (radius  $\sim 1$  fm). The quarks inside a hadron polarize the surrounding gluonic medium; the resulting gluon cloud around each quark provides it with a dynamically generated effective mass of about 300 MeV. In an ideal version of QCD, with massless quarks in the Lagrangian, this corresponds to spontaneous chiral symmetry breaking.

Confined hadronic matter exists in two distinct forms. At vanishing or low baryon number density, it consists largely of mesons, since the higher baryon mass reduces, through the Boltzmann factor, the baryonic thermodynamic weight. The interaction between mesons as well as that between mesons and baryons is resonance-dominated, i.e., it consists essentially of the abundant formation of multi-meson and meson-baryon resonances. Mesons appear to allow arbitrary overlap, so all (light quark) hadron states have the same characteristic size, with a radius of about 1 fm, independent of their mass. Nucleons, on the other hand, experience a short range repulsion in addition to a longer range attraction. The latter, the nuclear force, binds nucleons to nuclei, while the former leads to a nuclear volume linear in the total number of nucleons. This implies that nucleons, also with a hadronic radius of about 1 fm, have an effective hard core of about half that size. Both these forces are of non-resonant nature, so that the interaction in baryonic matter at low temperature and high density is thus quite different from that of mesonic matter. Nevertheless, in each case we have with increasing density, be it through “heating” or “compression”, a cluster formation which eventually leads to more quarks per hadronic volume than meaningful for a partitioning into color-neutral hadrons. In other words, increasing of the temperature  $T$  or the baryochemical potential  $\mu$  results eventually in color deconfinement. What happens in this transition?

On one hand, the color-neutral states are dissolved, producing a medium of color-charged constituents. The deconfinement transition is thus the QCD counterpart of the insulator-conductor transition of atomic physics. In addition, a sufficient increase of the temperature results eventually in a “melting” of the gluon cloud which surrounds the quarks inside a hadron. Hadronic matter thus shows two transitions, deconfinement and chiral symmetry restoration. Do these two phenomena necessarily coincide?

Rather general basic arguments [1] show that they either occur at the same point or if not, deconfinement precedes chiral symmetry restoration. A simplistic justification is given by the fact that any  $r$ -dependent confinement potential breaks chiral symmetry - a reflection of the quark at the potential wall does not flip its spin.

It is thus possible that quarks, when they become deconfined, still maintain their effective mass up to some higher temperature or density. Lattice calculations have shown that for vanishing baryon density, deconfinement and chiral symmetry restoration do in fact coincide, indicating that the deconfinement temperature is sufficient to melt the effective quark mass. For high baryon density at low temperature, this seems not likely, allowing a medium of massive quarks as an additional state of strongly interacting matter, in addition to hadronic matter and the plasma of deconfined massless quarks and gluons [2].

In any case, deconfinement does not result in a non-interacting medium; in section 3, we shall return in detail to the strong interaction in the QGP in the region above the onset of deconfinement. Here we only note that in particular the anti-triplet quark-quark interaction provides an attractive force, making possible the existence of diquarks as localized bound states. Such colored bosonic states can condense and therefore form a color superconductor as yet another state of strongly interacting matter. Putting everything together gives us the phase diagram shown in Fig. 2.

After this discussion, meant to show how the QGP fits into the general context of strong interaction thermodynamics, we shall now turn to its study in the low baryon number regime. Here extensive lattice calculations provide much information; moreover, this is also the region of relevance for RHIC and LHC experiments.

## 2. From Hadronic Matter to QGP

The simplest form of confined matter is an ideal gas of massless pions, whose pressure is given by

$$P_\pi = \frac{\pi^2}{90} 3 T^4 \approx \frac{1}{3} T^4, \quad (1)$$

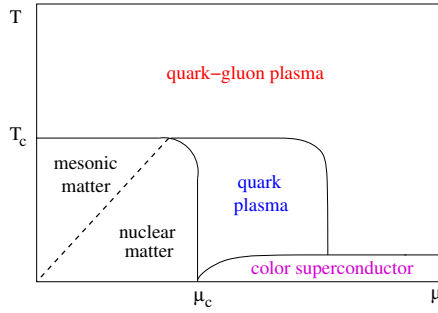


Figure 2: Speculative phase diagram of strongly interacting matter

taking into account the three possible pion charge states. For deconfined matter, we have as simplest case an ideal quark-gluon plasma (two massless quark flavors, two quark and two gluon spin orientations,  $q$  and  $\bar{q}$ ; three quark and eight gluon color degrees of freedom), giving

$$P_{QGP} = \frac{\pi^2}{90} \left\{ 2 \times 8 + \frac{7}{8} [2 \times 2 \times 2 \times 3] \right\} T^4 - B \simeq 4 T^4 - B \tag{2}$$

for the pressure, with  $B$  denoting the “bag pressure” exerted by the physical vacuum on the colored medium. The two forms of the pressure are compared in Fig. 3 (left).

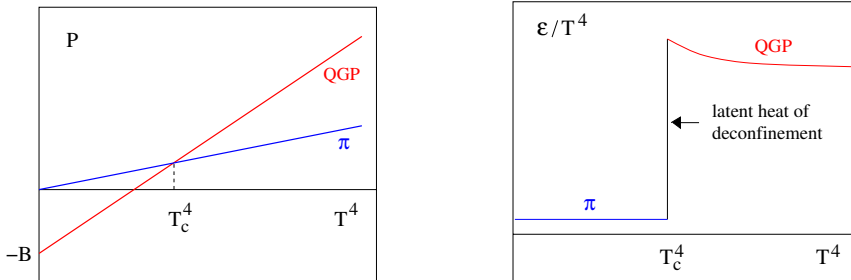


Figure 3: Pressures and energy densities for an ideal pion gas compared to an ideal QGP

Since nature always chooses the state of highest pressure (lowest free energy), this implies a phase transition from a pion gas at low  $T$  to a QGP at high  $T$ . The critical temperature is obtained from  $P_\pi = P_{QGP}$ , leading to  $T_c^4 \simeq 0.3 B$  and hence  $T_c \simeq 150$  MeV, if we use  $B^{1/4} \simeq 200$  MeV, as obtained from quarkonium spectroscopy. The corresponding energy densities become  $\epsilon_\pi \simeq T^4$  and  $\epsilon_{QGP} \simeq 12 T^4 + B$ , leading to the form shown in Fig. 3 (right).

The transition in this model is by construction of first order; at  $T_c$ , the energy density changes abruptly by the latent heat of deconfinement. For an ideal gas of massless constituents, the energy density and pressure are related by  $\epsilon = 3P$ ; in such a conformal world, the temperature is the only scale. In our case, the interaction measure

$$\Delta \equiv \frac{\epsilon - 3P}{T^4} = \frac{4B}{T^4} \tag{3}$$

is definitely not zero; even in this simplistic model, the QGP is non-conformal (with  $B$  as external scale) and strongly interacting for  $T_c \leq T < 3 - 4 T_c$ .

To turn from model to theory, we want to derive the thermodynamics obtained with QCD as dynamical input. The only possible *ab initio* calculations are based on the computer simulation [3] of the lattice regularization of QCD [4]. We therefore summarize the main results thus obtained in finite temperature lattice QCD.

In Fig. 4 we show the temperature behavior of the energy density for the cases of 2, 2+1 and 3 light quark flavors [5]; here 2+1 means one heavy and two light quarks. The sudden jump corresponding to the latent heat of deconfinement is quite evident. It is found to occur at a critical temperature of about 160 - 180 MeV, and the energy density at that point is around 0.5 to 1.0 GeV/fm<sup>3</sup>.

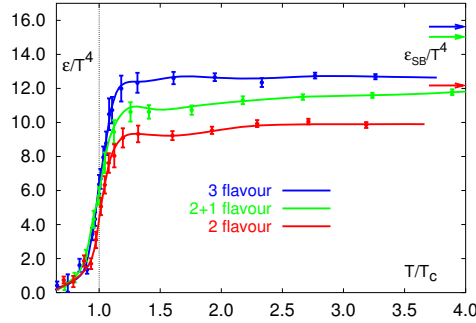


Figure 4: Lattice results for energy density vs. temperature in QCD thermodynamics [5]

To relate this “jump” more specifically to some form of critical behavior, we consider the corresponding order parameters for deconfinement and for chiral symmetry restoration. Such parameters signal the onset of the new phase.

For deconfinement, the order parameter is given by the average value of the Polyakov loop [6],

$$\langle L(T) \rangle \sim \exp\{-F_{Q\bar{Q}}(T)/T\}, \quad (4)$$

where  $F_{Q\bar{Q}}(T)$  denotes the free energy of a static  $Q\bar{Q}$  pair at infinite separation. In a confining medium, this diverges ( $F_{Q\bar{Q}}(r, T) \sim \sigma r$ ), while in a deconfined medium, screening prevents communication between the  $Q$  and the  $\bar{Q}$  beyond a certain distance, so that  $F_{Q\bar{Q}}$  remains finite. As a result,

$$L(T) \begin{cases} = 0 & T < T_L \text{ confinement} \\ \neq 0 & T > T_L \text{ deconfinement} \end{cases} \quad (5)$$

defines the deconfinement temperature  $T_L$ . Actually,  $\langle L(T) \rangle$  vanishes exactly for  $T < T_L$  only in the case of infinitely heavy quarks; for finite quark mass, the string binding  $Q$  and  $\bar{Q}$  breaks when the potential surpasses the spontaneous pair formation threshold. As a result,  $\langle L(T) \rangle$  is very small but finite for  $T < T_L$ .

The breaking of chiral symmetry is indicated by a finite value of the chiral condensate  $\chi(T) \equiv \langle \bar{\psi}\psi \rangle \sim M_q$ , which measures the dynamically generated (“constituent”) quark mass  $M_q$ , obtained for a Lagrangian with massless quarks. At high temperature, this mass melts, so that

$$\chi(T) \begin{cases} \neq 0 & T < T_\chi \text{ chiral symmetry broken} \\ = 0 & T > T_\chi \text{ chiral symmetry restored} \end{cases} \quad (6)$$

defines the chiral symmetry restoration temperature  $T_\chi$ . Here we have exact chiral symmetry only if the input quarks are massless; for finite quark mass, the symmetry remains explicitly broken, so that then  $\chi(T)$  only becomes very small for  $T > T_\chi$ .

Both  $\langle L(T) \rangle$  and  $\chi(T)$  have been studied extensively in finite temperature lattice QCD at vanishing overall baryon number. The corresponding susceptibilities (derivatives with respect to  $T$ ) peak sharply, defining  $T_L$  and  $T_\chi$ , and within errors, the two temperatures and hence the two phenomena (deconfinement and chiral symmetry restoration) coincide. The critical temperature for the resulting transition from hadronic matter to QGP, for two light quark flavors, is thus determined as  $T_c \simeq 175$  MeV.

It was clear from the very first finite temperature lattice studies that the deconfined medium in the region above  $T_c$  is very strongly interacting and thus quite far from an ideal plasma [7]. This is best seen from the interaction measure (the trace of energy-momentum tensor), defined as

$$\Delta = \frac{\epsilon - 3P}{T^4}. \quad (7)$$

For non-interacting massless constituents (the “conformal” limit),  $\Delta \equiv 0$ , so that the temperature is the only scale. The quarks and gluons of QCD are ideally massless, but the so-called trace anomaly violates conformality and introduces a dimensional scale. The behavior of  $\Delta(T)$  is shown in Fig. 2 both in pure gauge theory for different color groups and for full QCD with different flavor content.

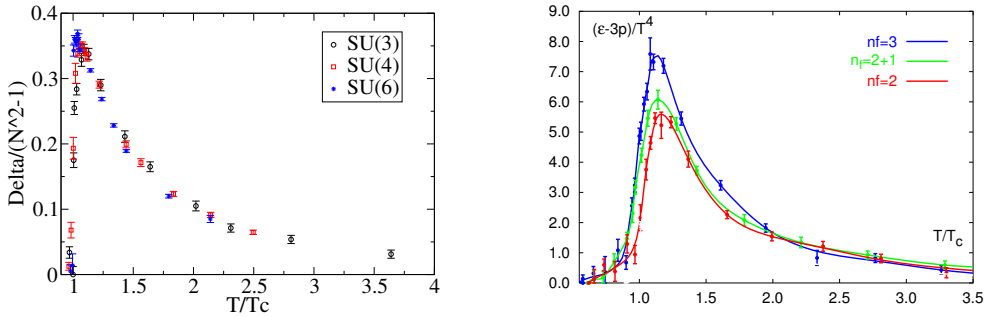


Figure 5: Interaction measure in  $SU(3)$  gauge theory (left) [8] and full QCD (right) [5].

### 3. The Strongly Interacting QGP

For sufficiently high temperature, asymptotic freedom is expected to result in an ideal QGP. How high does  $T$  have to be in order to allow some form of a weak coupling expansion (perturbation theory) to describe the approach to this limit?

Infrared divergences limit the perturbative expansion of QCD to  $O(g^5)$  [9]. The evaluation up to this order has been carried out [10] and is found to result in strongly oscillating and hence non-convergent behavior for  $T \leq 10 T_c$ .

This has led to considerable efforts to “repair” the difficulty, either by introducing non-perturbative scale effects to allow a systematic extension of perturbation theory beyond  $O(g^5)$  [11], or by regrouping sets of Feynman diagrams to expand around a ground state including screening effects (“resummed” perturbation theory [12], hard thermal loop approach [13]). In both cases, however, such weak-coupling methods cannot account for the behavior observed in lattice studies. This holds in particular for  $SU(3)$  gauge theory, where one has results for the continuum limit [14]; see Fig. 6 (left) for a comparison to HTL results. It is obvious that no weak-coupling approach can account for the critical behavior near  $T_c$  (the dip of  $\Delta(T)$  as  $T \rightarrow T_c$ ); but also the behavior in the region up to about  $5 T_c$ , with  $T^2 \Delta(T) \simeq \text{const.}$ , is not reproduced by the weak logarithmic form of perturbative studies.

The simplest non-perturbative approach, using the bag model form discussed above, does not fare any better. The bag pressure can be related to the gluon condensate  $G_0^2$  at  $T = 0$  [15], and using numerical estimates for the latter

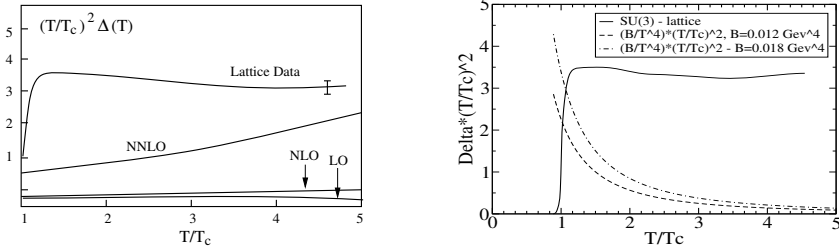


Figure 6: Interaction measure in  $SU(3)$  gauge theory [14], compared to (left) HTL resummed perturbation theory [13], and (right) bag model prediction [18].

[16], one obtains a  $\Delta(T)$  vanishing as  $T^{-4}$  and thus much too fast; moreover, the critical dip near  $T_c$  is also here not given (see Fig. 6 (right)).

We thus need to find a more detailed way to account for the non-perturbative behavior observed in the region  $T_c \leq T \leq 5T_c$ . One such possibility is given by the quasi-particle approach, first applied to  $SU(2)$  gauge theory [17]. The basic idea is that in the deconfined medium, gluons acquire an effective mass by polarizing the colored medium around them. In the critical region, as  $T \rightarrow T_c$  from above, the correlation length  $\xi(T)$  increases strongly or diverges, so that the gluon sees and polarizes more and more of the medium, increasing its mass. Outside the critical region, for  $T \geq 1.3T_c$ , as  $T$  increases further, the correlation length decreases essentially as  $T^{-1}$ ; now, however, the energy density  $\epsilon(T)$  increases (as  $T^4$ ), so that the effective mass  $M_g \sim \epsilon \times \xi \sim T$  also grows. These two competing modes lead to a growth of  $M_g$  for  $T \rightarrow T_c$  determined by the critical behavior and a conformal growth at high temperature. The dip in  $M_g(T)$  around  $T \sim 1.3T_c$  defines the transition from critical behavior to hot QGP. Applying this approach to the continuum limit of  $SU(3)$  gauge theory leads to excellent agreement, for the interaction measure as well as for the energy density, in both the critical region and in the hot QGP [18].

We conclude that in the region of main interest to us, for  $T_c \leq T \leq 3 - 5T_c$ , the QGP is strongly interacting. The interaction provides the constituents with a temperature-dependent effective mass, so that a considerable part of any energy input goes into mass rather than into kinetic energy. As a result of this, the behavior of the equation of state is highly non-ideal and in fact not unlike that just below  $T_c$ , where abundant resonance formation restricts the fraction of energy available for kinetic purposes.

#### 4. Probing the QGP

At sufficiently high temperatures and/or densities, strongly interacting matter thus becomes a plasma of deconfined, colored quarks and gluons. How can we probe the properties of this medium and study its behavior as function of temperature and density? This is a highly non-trivial problem, and I want to outline here briefly three ways of addressing it. We assume that we are given a macroscopic volume of deconfined strongly interacting matter and want to determine its state for different temperatures. This means that we shall consider equilibrium thermodynamics only, leaving out all the phenomena (collisions effects, time dependence, equilibration, flow) that make the analysis of actual nuclear collisions data so complex.

The given medium is by assumption hotter than its environment (the vacuum) and hence emits radiation. An outside observer will detect the emission of light hadrons; however, these cannot exist in the interior of the QGP and hence must be formed through hadronization at the cooler surface. Such radiation will therefore provide information about the hadronization stage of the QGP, but not about the pre-hadronic state in the interior. In the hot QGP itself, quark-gluon interactions and quark-antiquark annihilation produce real and virtual photons, respectively, and these will leave the medium without further strong interaction. They can thus provide information about the state of the medium when they were formed, i.e., about the hot QGP [19]. The difficulty is that they can be formed at all evolution stages of the medium, even in the hadronic phase, and so one has to find a way to identify hot thermal electromagnetic radiation. If this can be achieved, such radiation provides a thermometer for the medium.

Alternative tools are obtained by testing the medium with external probes. In particular, we can study the effect of the medium on quarkonia or on jets. Both will interact strongly in a deconfined medium and less or not at all with hadronic matter; thus they can provide information on the temperature and/or density of the QGP.

Quarkonia are bound states of heavy quark-antiquark pairs ( $c\bar{c}$ ,  $b\bar{b}$ ). They are much smaller than “light” hadrons ( $r_Q \ll r_h \sim 1$  fm) and much more tightly bound, with binding energies up to 0.5 to 1.0 GeV. Therefore they can survive in a QGP up to temperature above the deconfinement point and “melt” only when the color screening radius has dropped to quarkonium size [20]. Since the different quarkonium states have different sizes and binding energies, since will lead to a “sequential” suppression of quarkonia: first, the larger and more loosely bound excited states are dissolved, finally the small and tightly bound ground states. For charmonia, this is illustrated in Fig. 7, with  $\psi'$  and  $\chi_c$  melting followed eventually by that of the  $J/\psi$ . Such patterns can provide a spectral analysis of the QGP, similar to that obtained for the sun by solar spectra [21].

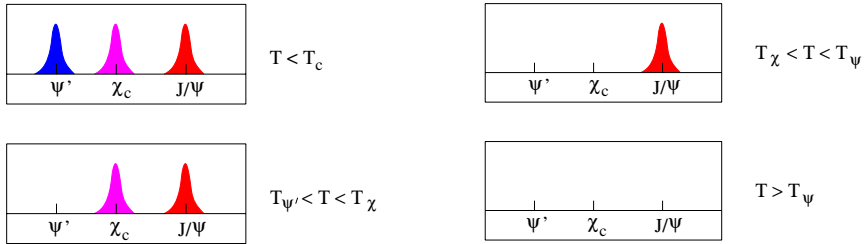


Figure 7: The spectral analysis of the QGP through charmonium states

Jets are fast partons (quarks or gluons) passing through the medium. They are colored and hence interact stronger with a QGP than with color-neutral hadronic matter. A substantial attenuation (“quenching”) of jets thus indicates the presence of a dense deconfined medium [22, 23].

We had called quarkonia and jets “external” probes. It is clear, however, that they have to be produced in the same collision which leads to the QGP candidate to be probed. They are, however, produced through very early hard interactions, which take place before the QGP is formed. We can then study the subsequent effect of the QGP on their behavior. Moreover, their initial production is to a large extent calculable by perturbative QCD, and it can be gauged in the study of  $pp$  and  $pA$  collisions, which presumably do not produce a QGP.

### 5. Three Questions to the LHC

The QGP predicted by statistical QCD is the ultimate state of matter to be studied in high energy nuclear collisions. This is a speculative endeavor, since it is not clear to what extent such collisions can produce something to be called matter. We therefore close our survey with three questions to the next generation of experiments which might help us in finding an answer to this fundamental enigma.

If an increase of collision energy indeed leads to the production of a hotter bubble of deconfined primordial matter, then this must expand more in order to reach the hadronization temperature, and hence the source size for hadron emission must become larger. In particular, it is expected to increase as a power of the hadron multiplicity, since this in turn grows with the initial energy density [24]. So far, from AGS to RHIC, the source size for hadron emission, as determined by Hanbury-Brown–Twiss (HBT) methods [25] used in astrophysics, has not shown a significant increase [26]. This “HBT-puzzle” has been accounted for in terms of the relative role of meson and baryon production [27], but at LHC energies, a clear increase of the source volume is predicted. Such an increase seems necessary in a model-independent way, if the concept of hot primordial fireball production in nuclear collisions is to make any sense.

We had noted that momentum spectra for real and virtual photons can in principle provide an internal thermometer of the QGP, with

$$(dN_\gamma/dk_T) \sim \exp\{-k_T/T\} \quad (8)$$

A recent analysis of RHIC  $Au-Au$  data at  $\sqrt{s} = 200$  GeV [28] has identified possible thermal photons, seen in a transverse momentum window between pion decay and prompt photon spectra. The corresponding temperature is with  $T = 221 \pm 19(\text{stat.}) \pm 19(\text{syst.})$  MeV above the hadronization value of about 175 MeV. If such thermal photons are indeed observable, the LHC should lead to much higher temperatures for electromagnetic radiation.

The last question addresses quarkonium production in nuclear collisions at the LHC. The  $J/\psi$  production rate in  $Au-Au$  collisions at RHIC is compatible with that for central collisions at the SPS, once cold nuclear matter effects are taken into account. The remaining survival rate of about 50 % is in accord with suppression of the higher excited states ( $\psi'$  and  $\chi_c$ ) and survival of the direct  $J/\psi$  [29]. The much higher energy density of the LHC should dissociate also the latter, leading to complete  $J/\psi$  suppression (modulo  $B$  decay and corona production). The expected survival pattern is illustrated in Fig. 8.

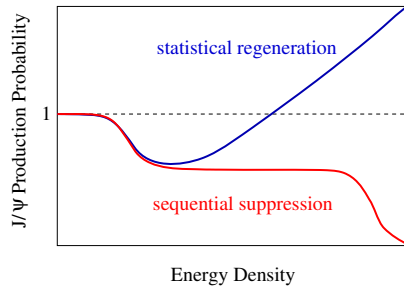


Figure 8: Sequential  $J/\psi$  suppression vs. statistical  $J/\psi$  regeneration

Here, however, an alternative scenario has been proposed [30] and much discussed. Charm production in nuclear collisions, as a hard process, increases with collision energy much faster than that of light quarks. At sufficiently high energy, the produced medium will therefore contain more charm quarks than present in a QGP at “chemical” equilibrium. If these charm and anticharm quarks combine at the hadronization point statistically to form charmonium states, this new combination mechanism should lead to a much enhanced  $J/\psi$  production rate, even if all primary (“direct”)  $J/\psi$ 's are dissociated. The two predictions, sequential suppression vs. statistical regeneration, thus present two really opposite patterns, and first LHC results should be able to distinguish between them. It should be emphasized that statistical recombination would on one hand provide clear evidence for the presence of a thermal medium, including even charm quarks. On the other, it presupposes a “new” statistical charmonium production mechanism, quite distinct from the hard production forms generally discussed, in which only  $c\bar{c}$  pairs at very short separation can bind to form a  $J/\psi$ .

## 6. Summary

We have seen that in strong interaction thermodynamics, there exists at vanishing baryon density a well-defined transition, in which

- color deconfinement sets in and chiral symmetry is restored,
- the energy density increases by the latent heat of deconfinement;
- the critical temperature  $T_c$  is about  $175 \pm 10$  MeV.

For  $T > T_c$ , the state of matter is a plasma of deconfined quarks and gluons, which can be probed by

- electromagnetic radiation,
- quarkonium spectra,
- jet quenching.

A more extensive survey of QCD thermodynamics will appear soon [31]. – Finally, we have given three rather model-independent bench-marks for a QGP study through very high energy nuclear collisions,



- does the source size (finally) increase with collision energy?
- does the thermal photon temperature increase with collision energy?
- does quarkonium production show sequential suppression or statistical regeneration?

The LHC should be able to provide some first answers to these question within the first year of heavy ion operation.

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