Hadronic Signals of Deconfinement at RHIC

Berndt Müller
Department of Physics, Duke University, Durham, NC 27708-0305, USA

This article reviews (soft) hadronic signals of deconfinement and chiral symmetry restoration in hot QCD matter in the light of the results from the first three years of the experimental program at the Relativistic Heavy Ion Collider.

I. INTRODUCTION

Quantum mechanics dictates that even "empty" space is not empty, but rather filled with quantum fluctuations of all possible kinds. In many contexts, such as in atomic physics, these vacuum fluctuations are subtle effects which can only be observed by precision experiments. In other situations, especially when interactions of sufficient strength are involved, the vacuum fluctuations can be of substantial magnitude and even "condense" into a nonvanishing expectation value of some quantum field. These vacuum condensates can act as a medium [1], which influences the properties of particles propagating through it.

An important example of such a vacuum condensate is the Higgs vacuum, which is introduced in the Standard Model of particle physics to generate the masses of quarks, leptons, and the gauge bosons of the weak interaction. The vacuum expectation value of the Higgs field, $\langle\phi\rangle = 246$ GeV, is uniquely determined in the Standard Model; the quark and lepton masses differ from one another only due to the different strength of the coupling of each fermion to the Higgs field. The quark masses receive additional contributions from the quark and gluon condensates in the QCD vacuum, which have been measured to be:

$$\langle\bar{q}q\rangle = (235\text{MeV})^3,$$
$$\langle g^2 G_{\mu\nu}G^{\mu\nu}\rangle = (840\text{MeV})^4. \qquad (1)$$

In fact, the contribution of the QCD vacuum condensates to the masses for the three light quark flavours u, d, s considerably exceed the mass believed to be generated by the Higgs field see Fig. 1).

If the vacuum acts as a medium and influences the properties of fundamental particles and their interactions, its properties can conceivably change. This idea has important implications in cosmology (inflation [2]) and in modern versions of the anthropic cosmological principle [3, 4], which invoke the concept of multiple universes with different phenomenological properties due to the realization of different vacuum states. The most promising candidate for the fundamental theory of spacetime, superstring theory, promises to sustain an almost unlimited number of vacuum states [5], lending some credence to this concept. These theoretical speculations call for an experimental verification of the existence of multiple vacuum states and the possibility of transitions among them under the influence of external constraints.

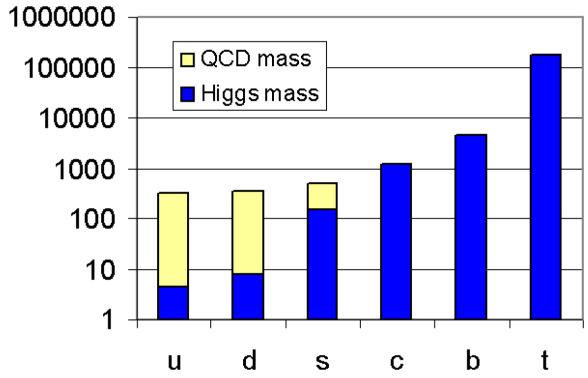

FIG. 1: Masses of the six quark flavors. The masses generated by electroweak symmetry breaking (current quark masses) are shown in dark blue; the additional masses of the light quark flavors generated by spontaneous chiral symmetry breaking in QCD (constituent quark masses) are shown in light yellow. Note the logarithmic mass scale.

Because of the energy scales involved, only the QCD vacuum is amenable to modification at energies accessible with present technologies. The Relativistic Heavy Ion Collider (RHIC) was built to explore these questions by creating conditions in the laboratory under which the structure of the QCD vacuum would change and the properties of particles and forces of the strong interactions would be modified. The beam energy of RHIC was chosen to be high enough to generate temperatures commensurate with the energy scale of the QCD vacuum condensates and thus to affect the structure of the QCD vacuum.

Numerical simulations of lattice QCD predict that a dramatic change in the QCD vacuum state occurs around the critical temperature $T_c \approx 160$ MeV, where the quark and gluon condensates melt, and the vacuum takes on a simpler structure. The part of the light quark masses that is induced by the quark condensate disappears above T_c and only the much smaller masses generated by the electroweak Higgs field remain. The degrees of freedom corresponding to independently propagating gluons, which are frozen in the normal QCD vacuum, are also liberated above T_c. This pattern is visible in Fig. 2, which shows the dramatic jump in the scaled energy density $\epsilon(T)/T^4$ at the critical temperature.

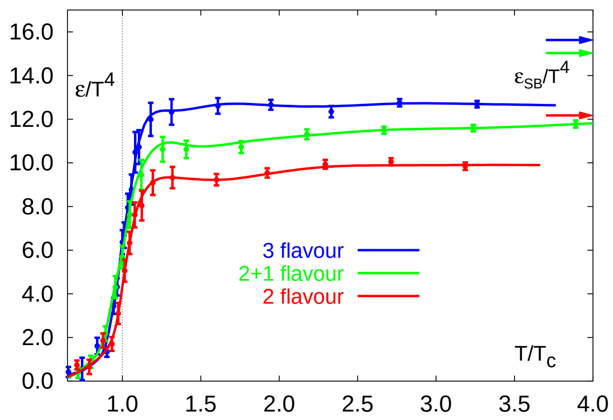

FIG. 2: Scaled energy density ϵ/T^4 for thermal lattice-QCD with two and three light quark flavors and for two light and one heavier flavor (from Karsch [43]).

II. HADRONIC PROBES OF DECONFINEMENT

Which observables allow us to confirm that the predicted transformation of the QCD vacuum actually occurs? Here we focus on signals involving light (u, d, s) quarks, the primary constituents of the medium. This restriction is not intended to minimize the importance of other observables, such as hard probes, covered in different contributions to this volume [7]. In fact, the strong quenching of the emission of hadrons at large transverse momenta and the existence of a strong (anisotropic) collective flow are important prerequisites for the apparent dominance of quark recombination at momenta of a few GeV/c, which will be discussed below.

Two effects caused by the disappearance of the vacuum condensates stand out: The melting of the quark condensate lowers the effective mass of the s-quark from about 500 MeV to less than 150 MeV, making it easy to create s-quark pairs in copious quantities [8–10]. The dissolution of the gluon condensate and the concomitant screening of the long-range color force by thermal gluons allows light quarks to propagate outside of hadronic bound states. As the hot quark-gluon plasma (QGP) cools, hadrons should form by recombinant emission of quark-antiquark pairs (mesons) or color-singlet triplets of quarks (baryons) from the surface of the hot medium [11], preserving the collective flow pattern of their partonic constituents [12].

A. Flavor equilibration

Let us consider these arguments more quantitatively. The expectation that the formation of a deconfined quark-gluon plasma results in a substantial enhancement of the production of hadrons carrying strangeness is based on a comparison of energy thresholds and rates. In a hadronic environment, the energy required to excite strangeness is about 700 MeV, as in the reaction $\pi\pi \to K\bar{K}$, or about $4T_c$. In the QGP, the threshold for the dominant reaction $gg \to s\bar{s}$ is about 300 MeV, less than $2T_c$. Hadronic reactions producing multiply strange hadrons, such as the Ξ and Ω hyperons, are thus predicted to be slow under thermal equilibrium conditions and not leading to hadrochemical equilibrium within the time available in a heavy-ion reaction [13]. On the other hand, flavor equilibrium among u, d, s quarks is predicted to develop within a QGP on the available time scales [8]. Hadrons containing any number of strange quarks should, therefore, be produced in equilibrium abundances when a QGP is formed.

The prediction of a strong enhancement of the production of multi-strange baryons and antibaryons is independent of the details of the hadronization model, although the quantitative predictions for individual hadron yields may differ slightly [10, 14–17]. Extensive studies of final state effects on the chemical composition of the emitted hadrons have concluded that the abundances of multi-strange baryons are only slightly affected by inelastic rescattering [18]. The consistency of this scenario was confirmed recently once more in a systematic study of baryon pair production based on model independent rate equations, which showed that a strong enhancement of multi-strange baryon production requires the assumption of flavor equilibrium already at the start of the hadronic phase evolution [19, 20].

It is important to distinguish the chemical equilibrium distribution attained in a deconfined plasma phase from the statistical distribution of hadrons produced in e^+e^- or p+p reactions[21–23]. In those reactions, the production of hadrons carrying strange quarks is systematically suppressed. Although this phenomenon can be explained as an effect of flavor conservation in a small volume of hadronic size, it is precisely the *disappearance* of this suppression in heavy-ion collisions, corresponding to the transition from a canonical ensemble to a macrocanonical one, which is the signature of the creation of locally deconfined QCD matter.

The observation of a strong enhancement of multi-strange baryon and antibaryon emission in Pb+Pb collisions by experiment WA97 [24] formed an important pillar of the claim that evidence for a new state of matter had been found at the CERN-SPS [25]. The results were confirmed and extended to lower beam energies in recent years by the NA57 collaboration [26]. The conclusion that a deconfined phase is formed under the conditions created by Pb+Pb collisions at CERN-SPS energies is also supported by a recent comprehensive analysis of the data for several relevant observables in the framework of a schematic, but realistic model of the space-time evolution of the collisions [27].

As impressive as these results are, the evidence remains indirect and relies on the argument that a hadrochemical equilibrium cannot be established by interactions among hadrons with confined quarks. The higher beam energies available at RHIC have made it possible to explore regions of phase space where the formation of hadrons from recombining quarks can be directly observed and the partonic origin of their collective flow pattern can be established. It is important to recognize that, in contrast to the strangeness enhancement itself, these phenomena were not quantitatively predicted before the observations were made. However, I will argue that the theoretical arguments, even if they were developed under the motivation of the first experimental results from RHIC, are sufficiently robust to provide compelling evidence for the formation of a deconfined quark phase, as more complete and statistically precise data are becoming available.

B. Quark recombination

Let us begin with a discussion of quark recombination (see Fig. 3a). The statistical coalescence of two or more particles to form a bound state is a widely used concept, which remains somewhat murky in many cases for two reasons. First, the formation of a bound state requires that either at least one of the two particles is off-shell or that another particle carries away the surplus energy. Second, the implicit reduction in the number of degrees of freedom during coalescence tends to conflict with the second law of thermodynamics, which dictates that the entropy cannot decrease. Fortunately, there are regions of phase space – hadrons emitted with highly epithermal momenta – where these concerns do not constitute serious obstacles. Hadrons emitted with a high momentum make an almost instantaneous transition from the dense medium into the surrounding vacuum. Another way to state this is to say that – observed from the fast moving hadron – the radiating surface appears highly Lorentz contracted and thus thin on a hadronic length scale. Energy conservation is therefore not an important constraint on the recombination process. Moreover, at high transverse momentum, the masses of quarks and hadrons can be neglected in first approximation with corrections of order $\mathcal{O}(m^2/p_T^2) \ll 1$. And finally, fast hadrons constitute only a very small fraction of all emitted particles, thus limiting possible violations of the second law of thermodynamics.

Denoting the quark phase-space distribution by $w_\alpha(p)$ (here α denotes color, flavor, and spin degrees of freedom) and employing light-cone coordinates along the direction of the emitted hadron, instantaneous recombination predicts the following meson and baryon spectra, respectively [28]:

$$E\frac{dN_M}{d^3P} = \int d\Sigma \frac{P\cdot u}{(2\pi)^3} \sum_{\alpha\beta} \int dx\, w_\alpha(xP^+)\bar{w}_\beta((1-x)P^+)|\phi_{\alpha\beta}^{(M)}(x)|^2, \tag{2}$$

$$E\frac{dN_B}{d^3P} = \int d\Sigma \frac{P\cdot u}{(2\pi)^3} \sum_{\alpha\beta\gamma} \int dx\, dx'\, w_\alpha(xP^+)w_\beta(x'P^+)w_\gamma((1-x-x')P^+)|\phi_{\alpha\beta\gamma}^{(B)}(x,x')|^2. \tag{3}$$

Here Σ stands for the freeze-out hypersurface, u^μ for the collective flow velocity of the partonic matter, and capital letters P denote hadron coordinates. For a thermal parton distribution, $w(p) = \exp(p\cdot u/T)$, the partonic factors in both equations combine to a single thermal distribution for the hadron, $w(P^+)$. In this case the integrations over the momentum fractions x, x' become trivial, and the end result is independent of the details of the hadron wavefunction ϕ. The prediction is that hadrons are emitted in thermal and chemical equilibrium abundances and acquire the collective radial flow velocity of the quarks. In fact, one can argue that the recombination mechanism provides the justification for applying the statistical model to describe ratios of hadron yields at up to momenta of several GeV/c.

In the absence of a thermal medium, energetic hadrons are created by the fragmentation of a fast quark or gluons (see Fig. 3b). This mechanism predicts a hadron spectrum of the form

$$E\frac{dN_h}{d^3P} = \int d\Sigma \frac{P\cdot u}{(2\pi)^3} \sum_{\alpha} \int dz\, z^{-3}\, w_\alpha(P^+/z) D_{\alpha\to h}(z), \tag{4}$$

where $z < 1$ denotes the momentum fraction carried by the hadron. It is easily shown that for large values of the momentum P^+ recombination always wins over fragmentation in the case of an exponentially falling momentum distribution, because $w(P^+/z)$ falls off faster than $w(P^+)$. On the other hand, fragmentation asymptotically dominates over recombination, when the parton distribution obeys a power law as it is always the case in QCD at sufficiently large momentum. In relativistic heavy ion collisions the power law tail is suppressed (in central Au+Au collisions at RHIC by about a factor 5) due to the energy loss of fast partons dur-

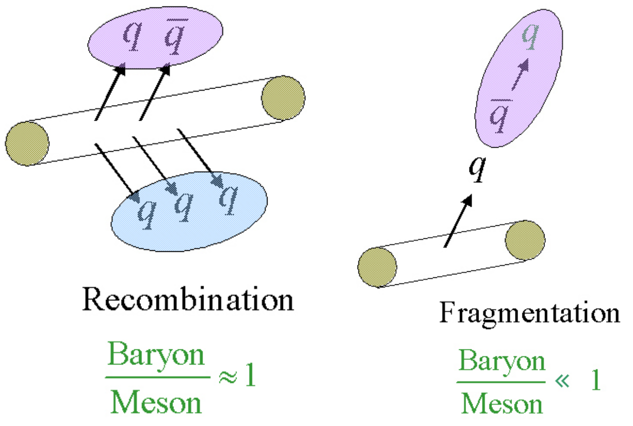

FIG. 3: (a) Left panel: Mesons and baryons formation by recombination of quarks and/or antiquarks from a thermal, deconfined medium predicts a baryon/meson ratio of order unity. (b) Right panel: Hadron production by fragmentation of an energetic parton favors meson over baryon production.

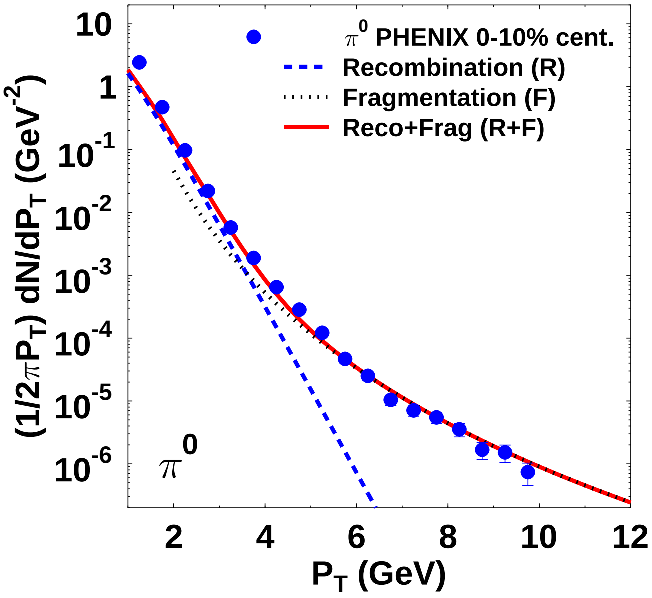

FIG. 4: Spectrum of pions emitted in central Au+Au collisions at RHIC for $\sqrt{s_{\mathrm{NN}}} = 200$ GeV. The contributions from recombination and fragmentation are shown separately, together with the data from the PHENIX collaboration (from [28]).

ing their passage through the hot, dense matter. This mechanism delays the dominance of fragmentation over recombination to rather large transverse momenta, about 4 GeV/c for mesons and 6 GeV/c for baryons. This is illustrated in Fig. 4, which shows the two contributions to the pion spectrum in central Au+Au collisions at RHIC separately, together with the PHENIX data.

C. Collective flow of quarks

When the radial flow velocity of the QGP exhibits an azimuthal anisotropy, as in noncentral collisions of two nuclei, this anisotropy becomes imprinted on the hadrons formed by quark recombination. If the quark flow anisotropy has the form $[1+v_2(p)\cos 2\varphi_p]$, where φ_p is the emission angle with respect to the collision plane and $v_2(p)$ the so-called elliptic flow velocity, the elliptic flow parameter for a meson or baryon is (in the usual case where $v_2 \ll 1$):

$$\begin{aligned} v_2^{(M)}(P) &= v_2(xP) + v_2((1-x)P), \\ v_2^{(B)}(P) &= v_2(xP) + v_2(x'P) + v_2((1-x-x')P). \end{aligned} \tag{5}$$

In the limiting case of equal momentum fractions carried by the constituent quarks these relations simplify to [12, 28, 29]

$$\begin{aligned} v_2^{(M)}(P) &\approx 2v_2(P/2), \\ v_2^{(B)}(P) &\approx 3v_2(P/3). \end{aligned} \tag{6}$$

In the hydrodynamic regime, where the flow anisotropy grows roughly linearly with the momentum, $v_2(p) = \alpha p$, the scaling factors cancel, and mesons and baryons exhibit an equal v_2 at the same momentum. At large transverse momentum, however, where the anisotropy is limited by transparency effects, this is no longer the case, and a collective anisotropic flow carried by quarks can be distinguished from one carried by a hadronic fluid. Final state effects on the elliptic flow in the hadronic phase have been shown to be small [30].

D. Unified descriptions

These considerations based on general principles can be extended in several ways to cover a wider range of hadron momenta [31]. By choosing the medium rest frame instead of the light-cone frame of the emitted hadron, one can obtain a more traditional coalescence model description of the hadron spectrum down to $p_T = 0$ [32]. The price one pays for this extended coverage is the need to make specific assumptions about the internal wavefunction of the emitted hadrons, which are by necessity simplistic. A virtue of this approach is that one can incorporate final-state decays of unstable hadrons, such as the ρ-meson, which help to resolve the entropy problem and improve the description of pion spectra, especially at small momenta [33]. Effects due to hadron masses can also be treated more consistently in such models.

Another extension concerns the treatment of fragmentation processes. In the presence of a medium, a fast parton may form a hadron by picking up one or more constituents from the medium. This process has been called fast-slow recombination [32]. Depending on the hadron wavefunction, it may contribute to hadron emission up to quite large transverse momenta extending into the $6-8$ GeV/c range. Alternatively, one can consider the process of hadron formation via fragmentation also as a recombination process acting on a primary distribution of fragmented shower partons [34, 35]. This approach gives a microscopic description of the fragmentation functions $D_{\alpha\to h}(z)$, allowing one to extend the calculation consistently, albeit model dependently, to hadron production by recombination involving partons from a thermal medium [36]. This model predicts a large baryon-to-meson ratio in the p_T-range of $3-8$ GeV/c, where the mixed recombination of thermal and shower partons dominates. This approach is also aplicable to hadron production in p+p and d+Au collisions, where it provides an explanation for the differences in the Cronin effect observed for mesons and baryons [37].

E. Principal signatures

These considerations suggest several critical signatures of the formation of a thermal quark-gluon plasma accompanied by a structural change in the QCD vacuum:

- Hadrons containing any number of strange quarks should be produced in near equilibrium abundances;
- Hadrons created by recombination of moderately energetic quarks should form the dominant mode of hadron emission at low and intermediate momenta;
- Baryons should be produced as abundantly as mesons at intermediate transverse momenta;
- The collective flow patterns of the emitted hadrons should reflect the collective flow of their constituent quarks.

As we will discuss in the next section, these characteristic phenomena have been observed in Au+Au collisions at RHIC.

III. RESULTS FROM RHIC

Fig. 5 shows that hadrons containing any number of s-quarks are produced according to the expectation that a chemically equilibrated quark-gluon plasma converts into hadrons maintaining flavor equilibrium near the critical temperature. The value deduced from the data ($T_{\rm ch}=177$ MeV [38]) is equal to T_c within the present theoretical uncertainties. There exist significant deviations from the predicted equilibrium abundances for some short-lived hadronic resonances, such as the $\Delta(1232)$, K^*, and the $\Lambda(1520)$. These deviations can be qualitatively understood as final state effects due to medium modifications of the resonance spectral function or absorption and regeneration effects [39].

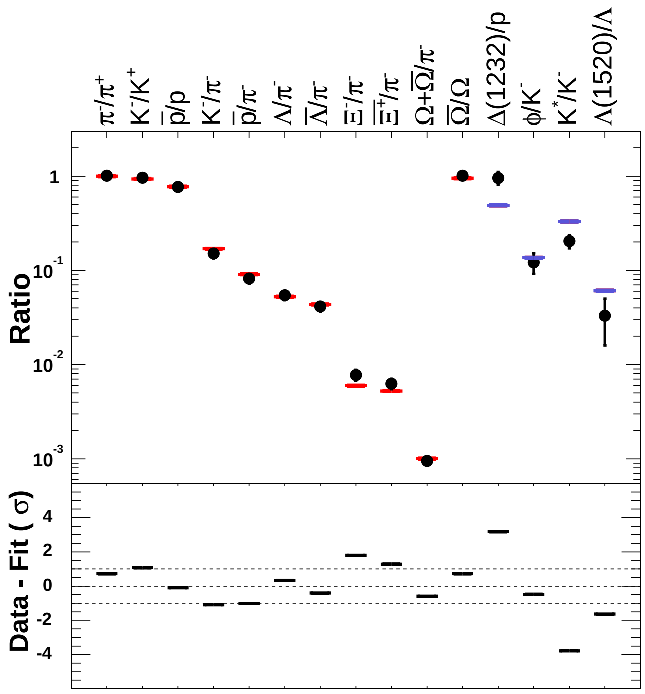

FIG. 5: Comparison of the ratios of observed yields of various hadrons at midrapidity in Au+Au collisions at RHIC with the ratios predicted in a thermal and chemical equilibrium model (from STAR).

The apparent lack of a nuclear suppression of protons and Λ-hyperons in Au+Au collisions in the transverse momentum range $2-4$ GeV/c, where pions are already suppressed by a factor $4-5$, came as a surprise. The measurements have been extended to other identified mesons and baryons, including the ϕ and the Ξ (see Fig. 6), confirming the baryon-meson dichotomy which is a generic feature of the recombination mechanism. The observed onset of nuclear suppression for the baryons between 4 and 6 GeV/c is in excellent agreement with the predictions by the Duke and Texas A&M models (see Fig. 7) and reflects the transition from a recombination dominated to s fragmentation dominated hadronization regime. The fall-off of the baryon/meson ratio at small p_T (see Fig. 7a) is due to mass effects.

Measurements of the flow patterns of the emitted hadrons provide essential additional evidence for the hypothesis that they are created directly from a phase of unconfined quarks and antiquarks. The experiment makes use of the fact that the region of heated vacuum in semiperipheral collisions between two nuclei is almond-shaped and thus sustains anisotropic pressure gradients. In the hydrodynamic limit this leads to an anisotropic expansion pattern called "elliptic flow" [40], which is characterized by the parameter v_2. As already mentioned, the flow anisotropy is an increasing function of the momentum of the emitted particle in the hydrodynamic regime and saturates when transparency effects set in at large transverse momenta.

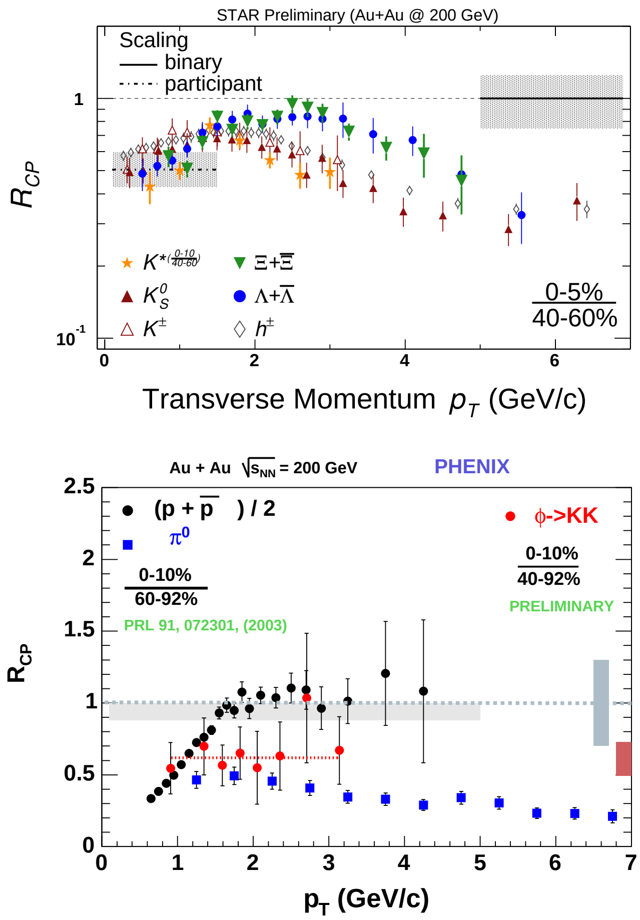

FIG. 6: Nuclear suppression observed in central (compared with peripheral) Au+Au collisions for various hadrons. (a) Upper panel: STAR data; (b) lower panel: PHENIX data.

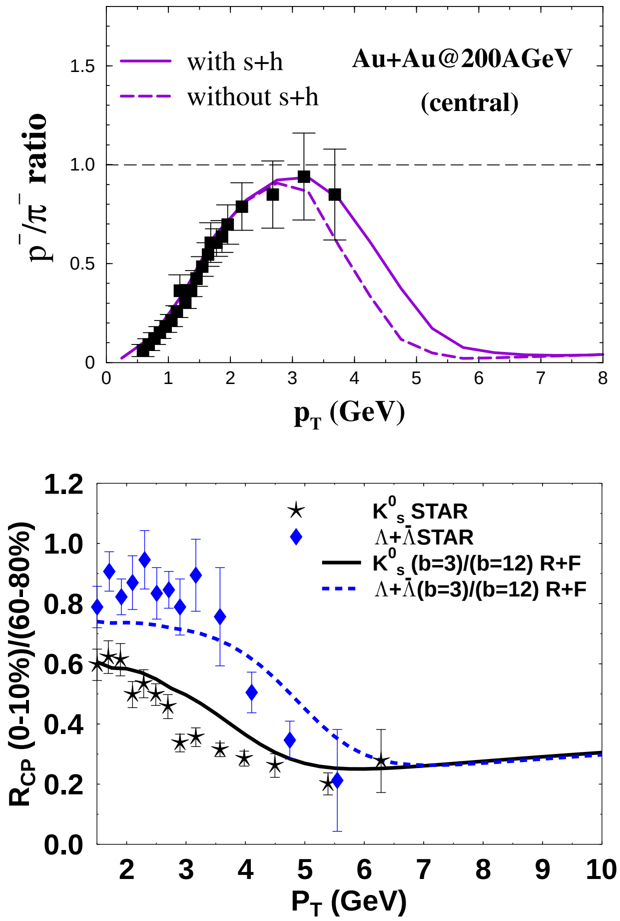

FIG. 7: Baryon enhancement as function of transverse momentum predicted by two different recombination models, in comparison with RHIC data. (a) Upper panel: Texas A&M model for p/π. (b) Lower panel: Nuclear suppression of Λ-hyperons and kaons in the Duke model.

The RHIC experiments have confirmed this expected behavior, but found that it varies sustantially from one hadron species to another (see Fig. 8a). However, when the constituent rescaling (6) is applied to these data, the flow patterns are found to collapse onto a common line (see Fig. 8b), suggesting that one observes the flow pattern of individual quarks, which coalesce into hadrons only at the end of the expansion [41]. The onset of saturation at a quark momentum $p_T \approx 1$ GeV/c is consistent with the transition between recombination and fragmentation observed in the baryon/meson ratios. The different behavior of mesons and baryons is characteristic of the quark recombination mechanism and reproduced in all model calculations [32, 42] (see Fig. 9).

IV. SUMMARY AND OUTLOOK

Experiments with relativistic heavy ions are allowing us, for the first time, to verify that a vacuum state can be modified under extreme conditions, causing dramatic changes in the properties of fundamental particles (quarks) and the strong forces among them. Hadrons emerging by recombination of the constituents of the hot QCD medium provide evidence for its flavor chemical equilibration and the deconfinement of (constituent) quarks. While flavor equilibration was already observed in experiments at the CERN-SPS, the characteristic momentum dependence of the baryon/meson ratios and the anisotropic collective flow patterns of hadrons have been observed at RHIC for the first time. Both features are model independent predictions of the quark recombination mechanism in a kinematic range where general arguments about the dominance of recombination apply if, and only if, a thermally and chemically equilibrated quark plasma hadronizes.

The present data and their theoretical descriptions leave a number of important questions unaddressed. One concerns two-body correlations among energetic hadrons. Recombination from a purely thermal medium only allows for quantum (HBT) correlations among hadrons,

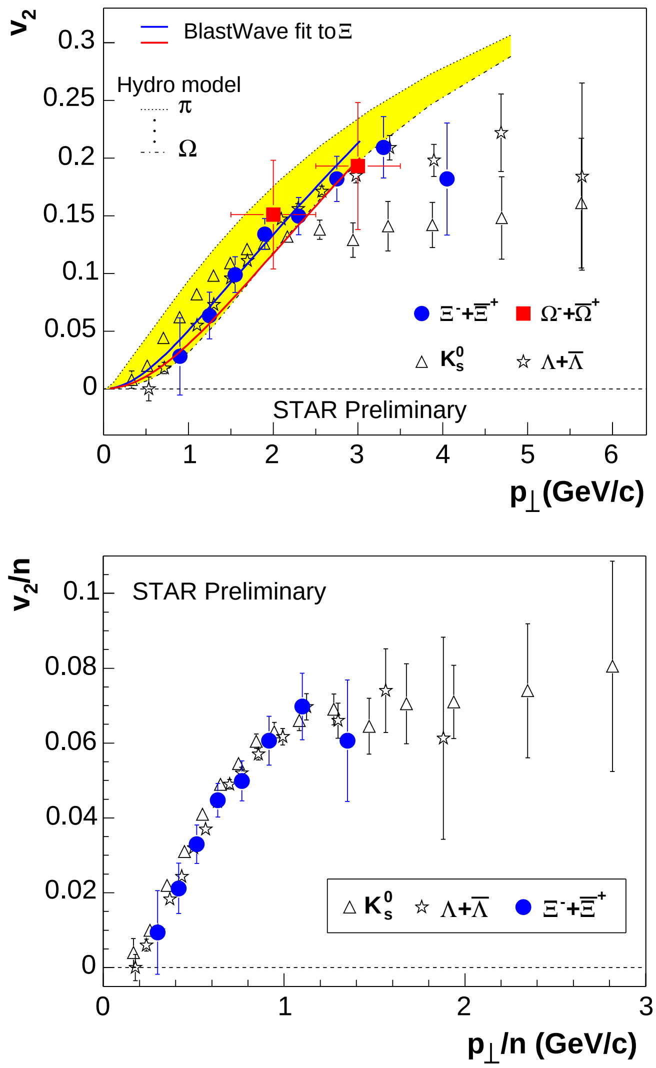

FIG. 8: (a) Upper panel: Elliptic flow parameter v_2 measured for various strange hadrons as function of transverse momentum p_T (STAR data). (b) Lower panel: Same data, but plotted in the scaled variables v_2/n and p_T/n, where $n = 2, 3$ is the number of constituent quarks in a hadron.

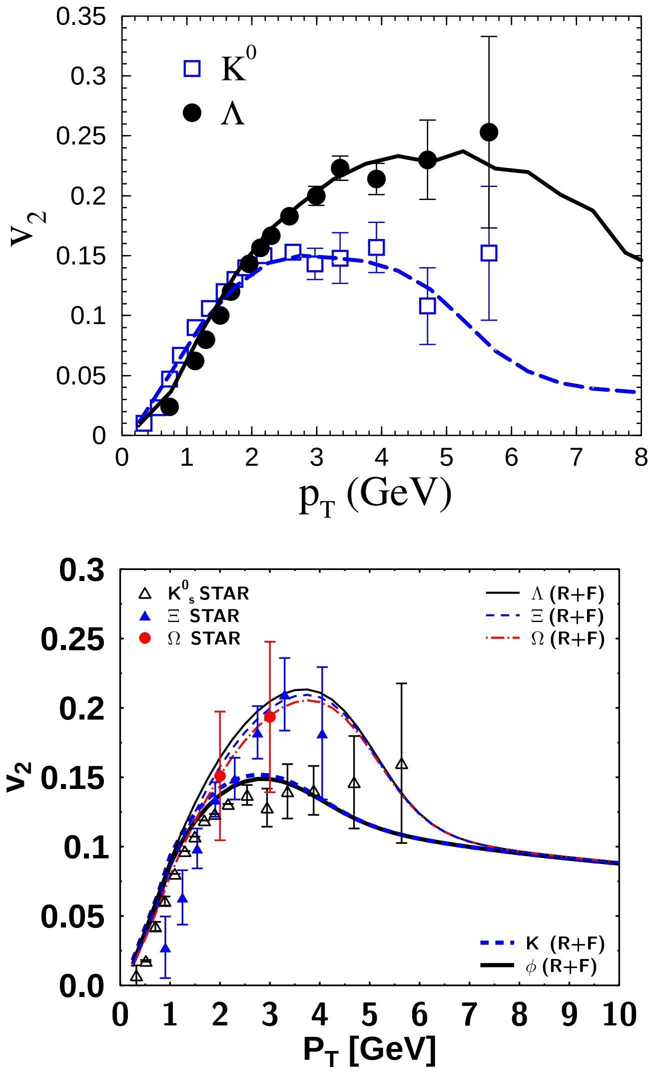

FIG. 9: Elliptic flow parameter v_2 as function of transverse momentum p_T together with the STAR data. (a) Upper panel: Kaons and Λ hyperons (Greco et al. [32]). (b) Lower panel: Kaons, Ξ and Ω hyperons (Nonaka et al. [42]).

but there is clear evidence in the data for the persistence of jet-like correlations into the kinematic regime where the recombination mechanism is thought to dominate. One explanation could be that these correlations are vestiges of the recombination of thermal and shower partons. Another explanation could be that the assumption of a completely thermal medium is a simplification, and that two-parton correlations become increasingly important with increasing p_T. These two explanations are likely just two faces of the same mechanism, because correlations among partons are probably caused by interactions of an energetic parton with constituents of the medium.

Another important puzzle is the apparent absence of gluonic degrees of freedom in the recombination mechanism. Does this indicate that gluons disappear as independent degrees of freedom as the plasma approaches the critical temperature from above? The precipitous drop in the Debye screening mass of a thermal plasma found in (quenched) lattice-QCD simulations near, but above T_c (see Fig. 8 of [43]) would suggest that this is, indeed, the case. One way to explain the absence of gluons is to argue that they fragment into quark-antiquark pairs prior to recombination [35].

The presence of strong dynamical correlations (clustering) among the constituents of the medium prior to hadronization, as suggested by theoretical considerations [44], would certainly serve to enhance the probability of recombination. This would be similar as the preformation of four-nucleon clusters facilitating the α-particle decay of heavy nuclei [45].

Acknowledgments

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[1] T. D. Lee, *Trans. N.Y. Acad. Sci.* (Ser.2) **40**, 111 (1980).
[2] A. H. Guth, *Phys. Rev. D* **23**, 347 (1981); A. D. Linde, *Phys. Lett.* **108B**, 389 (1982); A. Albrecht and P. J. Steinhardt, *Phys. Rev. Lett.* **48**, 1220 (1982). For a concise review of inflationary cosmology see, e. g. A. H. Guth, *Inflation and Eternal Inflation*, arXiv:astro-ph/0002156.
[3] M. J. Rees, *Astrophys. & Space Sci.* **285**, 95 (2003);
[4] P. C. W. Davies, *Mod. Phys. Lett. A* **19**, 727 (2004).
[5] L. Susskind, *The anthropic landscape of string theory*, arXiv:hep-th/0302219.
[6] F. Karsch, *Nucl. Phys. A* **698**, 199 (2002).
[7] Evidence for the formation of a quark-gluon plasma in Au+Au collisions at RHIC from various other signature observables is discussed in various other articles in this volume.
[8] J. Rafelski and B. Müller, *Phys. Rev. Lett.* **48**, 1066 (1982).
[9] J. Rafelski, *Phys. Rept.* **88**, 331 (1982).
[10] P. Koch, B. Müller and J. Rafelski, *Phys. Rept.* **142**, 167 (1986).
[11] A. P. Vischer, P. J. Siemens and J. C. Parikh, *Phys. Rev. D* **43**, 271 (1991).
[12] S. A. Voloshin, *Nucl. Phys. A* **715**, 379 (2003).
[13] P. Koch and J. Rafelski, *Nucl. Phys. A* **444**, 678 (1985).
[14] T. S. Biro and J. Zimanyi, *Nucl. Phys. A* **395**, 525 (1983).
[15] H. W. Barz, B. L. Friman, J. Knoll and H. Schulz, *Nucl. Phys. A* **484**, 661 (1988).
[16] T. S. Biro, P. Levai and J. Zimanyi, *Phys. Lett. B* **347**, 6 (1995).
[17] J. Rafelski and J. Letessier, *Phys. Rev. Lett.* **85**, 4695 (2000).
[18] S. A. Bass and A. Dumitru, *Phys. Rev. C* **61**, 064909 (2000).
[19] J. Kapusta and I. Shovkovy, *Phys. Rev. C* **68**, 014901 (2003).
[20] P. Huovinen and J. I. Kapusta, *Phys. Rev. C* **69**, 014902 (2004).
[21] R. Hagedorn, *Nuovo Cim. Suppl.* **3**, 147 (1965); R. Hagedorn and J. Ranft, *ibid.* **6**, 169 (1968).
[22] F. Becattini, A. Giovannini and S. Lupia, *Z. Phys. C* **72**, 491 (1996).
[23] F. Becattini and U. W. Heinz, *Z. Phys. C* **76**, 269 (1997).
[24] E. Andersen *et al.* [WA97 Collaboration], *Phys. Lett. B* **449**, 401 (1999).
[25] U. W. Heinz and M. Jacob, *Evidence for a new state of matter: An assessment of the results from the CERN lead beam programme*, arXiv:nucl-th/0002042.
[26] L. Sandor *et al.* [NA57 Collaboration], *J. Phys. G* **30** (2004) S129.
[27] T. Renk, *A comprehensive description of multiple observables in heavy-ion collisions at SPS*, arXiv:hep-ph/0403239.
[28] R. J. Fries, B. Müller, C. Nonaka and S. A. Bass, *Phys. Rev. Lett.* **90**, 202303 (2003); *Phys. Rev. C* **68**, 044902 (2003).
[29] D. Molnar and S. A. Voloshin, *Phys. Rev. Lett.* **91**, 092301 (2003).
[30] Z. W. Lin and C. M. Ko, *Phys. Rev. C* **65**, 034904 (2002).
[31] An up-to-date survey of the various approaches to quark recombination can be found in: R. J. Fries, *Recombination models*, arXiv:nucl-th/0403036.
[32] V. Greco, C. M. Ko and P. Levai, *Phys. Rev. Lett.* **90**, 202302 (2003); *Phys. Rev. C* **68**, 034904 (2003).
[33] V. Greco and C. M. Ko, *Effect of resonance decays on hadron elliptic flows*, arXiv:nucl-th/0402020.
[34] K. P. Das and R. C. Hwa, *Phys. Lett. B* **68**, 459 (1977); Erratum *ibid.* **73**, 504 (1978).
[35] R. C. Hwa and C. B. Yang, *Phys. Rev. C* **66**, 025205 (2002); *Phys. Rev. C* **67**, 034902 (2003).
[36] R. C. Hwa and C. B. Yang, *Recombination of shower partons at high p_T in heavy-ion collisions*, arXiv:nucl-th/0401001.
[37] R. C. Hwa and C. B. Yang, *Final-state interaction as the origin of the Cronin effect*, arXiv:nucl-th/0403001.
[38] P. Braun-Munzinger, K. Redlich and J. Stachel, *Particle production in heavy ion collisions*, arXiv:nucl-th/0304013, in: *Quark Gluon Plasma 3*, eds. R. C. Hwa and X.-N. Wang (World Scientific, Singapore, 2004).
[39] H. B. Zhang [STAR Collaboration], *Δ, K^* and ρ resonance production and their probing of freeze-out dynamics at RHIC*, arXiv:nucl-ex/0403010.
[40] J. Y. Ollitrault, *Phys. Rev. D* **46**, 229 (1992).
[41] J. Adams *et al.* [STAR Collaboration], *Phys. Rev. Lett.* **92**, 052302 (2004).
[42] C. Nonaka, R. J. Fries and S. A. Bass, *Phys. Lett. B* **583**, 73 (2004); C. Nonaka, B. Müller, M. Asakawa, S. A. Bass and R. J. Fries, *Phys. Rev. C* **69**, 031902 (2004)
[43] O. Kaczmarek, F. Karsch, E. Laermann and M. Lütgemeier, *Phys. Rev. D* **62**, 034021 (2000).
[44] E. V. Shuryak and I. Zahed, *Towards a theory of binary bound states in the quark gluon plasma*, arXiv:hep-ph/0403127.
[45] There exists an extensive literature addressing this problem. A representative example is: S. M. Lenzi, et al., *Phys. Rev. C* **48**, 1463 (1993).