THE QCD PHASE TRANSITION:
FROM THE MICROSCOPIC MECHANISM TO SIGNALS

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This talk consists of two very different parts: the first one deals with non-perturbative QCD and physics of the chiral restoration, the second with rather low-key (and still unfinished) work aiming at obtaining EOS and other properties of hot/dense hadronic matter from data on heavy ion collisions. The microscopic mechanism for chiral restoration phase transition is a transition from randomly placed tunneling events (instantons) at low T to a set of strongly correlated tunneling-anti-tunneling events (known as instanton - anti-instanton molecules) at high T. Many features of the transition can be explained in this simple picture, especially the critical line and its dependence on quark masses. This scenario predicts qualitative change of the basic quark-quark interactions around the phase transition line, with some states (such as pion-sigma ones) probably surviving even at $T > T_c$. In the second half of the talk we discuss experimental data on collective flow in heavy ion collision, its hydro-based description and relation to equation of state (EOS). A distinct feature of the QCD phase transition region is high degree of “softness”, (small ratio pressure/energy density). We present some preliminary results indicated that it is indeed needed to explain the radial flow at SPS energies.
1. New mechanism for the chiral phase transitions

The QCD-like theories with variable number of colors \( N_c \) and (light) flavors \( N_f \) have very rich phase structure, which only now starts emerging from theory and lattice simulations. It shows how naive are many textbook-style explanations, considering “overlapping hadrons” and “percolating quarks”. However, in all cases transitions happen at rather dilute stage of the hadronic phase\(^1\). Naive geometric ideas cannot explain why light fermions play such an important role. In particular, the critical temperature \( T_c \) for pure gauge simulations \( (N_f = 0) \) and those with dynamical quarks \( (N_f = 2-3) \) differ by almost factor 2, being \( T_c \approx 260\,\text{MeV} \) and \( T_c \approx 150\,\text{MeV} \) respectively. Furthermore, at \( T \approx T_c \) many physical quantities are very sensitive to such little details as the mass value of the strange quark. (That is why lattice results are still not quite definite about the order of the transition in the real world.) Furthermore, on the phase diagram as a function of quark masses there are 2 distinct region of the 1-st order transitions: the one at large masses is referred to as “deconfinement” and the one at low masses “chiral restoration\(^2\)."

An important aspects of the chiral restoration is the role of the U(1) chiral symmetry, see \[1\]. The standard arguments suggest that the \( \sigma \) meson (the \( SU(N_f) \)A chiral partner of the pion) becomes massless at \( T \rightarrow T_c \). The question is what happens with their U(1) chiral partners, \( \eta' \) and isovector scalar (with the old name \( \delta \) and new one \( a_0 \)) do. It remains unknown what happens with their masses, and some lattice data suggest those become very light as well.

Apart of (i) increasing the temperature \( T \), there are other things one can do to a QCD vacuum, and see when (if at all) the quark condensate disappear and chiral symmetry is restored. For example, one can (ii) increase the baryon density \( n_b \), (iii) increase the number of quarks \( N_f \) in the theory, or (iv) apply the magnetic field. There are indications that the mechanism I am going to discuss work in all cases (i-iii). However, in contrast to superconductivity, magnetic field does not destroy the quark condensate at all (see \[2\] and references therein).

By now there is large amount of evidence that in vacuum the chiral breaking and other phenomena related with hadrons made of light quarks are dominated by instantons, see recent review \[3\]. The instanton-based models explain multiple correlation functions and hadronic spectroscopy, and are also directly supported by lattice studies, e.g. \[8\].

Now we are ready to discuss the mechanism underlying chiral restoration. It is structural rearrangement of ensemble of instantons, from relatively random liquid at low \( T \) to a gas of instanton anti-instanton “molecules\(^3\). In the case of high-\( T \) transition one can explain what happens at \( T \approx T_c \) rather simply\(^4\). Recall that finite \( T \) field theory can be described in Euclidean space-time by simply imposing finite periodic box in time, with

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1. Pure \( SU(3) \) gluodynamics is an especially good example: at \( T_c \approx 260\,\text{MeV} \) the density of glueballs is negligible since even the lightest one has a mass of about 1.7 GeV!
2. These different names are partially misleading. We remind that for light fermions deconfinement has no well defined order parameter; nevertheless in both phase transitions the high-\( T \) phase is close to perturbative quark-gluon plasma, at least this is what EOS data tell us.
3. Note a similarity to Kosterlitz-Thouless transition in \( O(2) \) spin model in 2 dimensions: there are paired topological objects (vortices) in one phase and random liquid in another.
4. Unfortunately, in rather technical terms. Simpler explanation is as follows: tunneling and antitunneling events at \( T > T_c \) happen at about the same place.
the Matsubara period $\tau = 1/T$. A rising $T$ means shrinking box size, and it turns out the transition happens exactly when one $\Pi$ molecule fits into the box. In a series of recent numerical simulations of the instanton ensembles (see review [3]) this phenomenon is very clearly observed. Not only we get chiral restoration at the right temperature, the order of the phase transition and its dependence on quark masses agrees with lattice data. Transition looks like the second order for $N_f = 2$ massless flavors, but turns to a weak first order one for QCD with physical masses. Furthermore, the thermodynamic parameters, the spectra of the Dirac operator, the $T$-dependence of the quark condensate and various susceptibilities, the screening masses and even the critical indices are consistent with available lattice data.

Few words about QCD with larger $N_f$. Fig.4 show that at $N_f = 5$ instantons can no longer break the chiral symmetry\textsuperscript{5} even at $T=0$ (provided quarks are light enough). On the other hand, if instantons cannot break chiral symmetry, other effects like confinement or even one gluon exchange can do it, and according to [4] this happens for $N_f \simeq 11$. Existing estimates\textsuperscript{2} show that the quark condensate in the domain $N_f = 5 - 11$ is small, with a value from about 1/10 of the QCD value to the exponentially vanishing one at the upper end. Lattice measurements in this domain so far does not see it, and report chirally symmetric ground state at 8,12 and 16 flavors.

The last topic is hadron modification near the phase transition line. Their masses are expected to change as quark condensate decreases, and the most radical idea is the Brown-Rho scaling, suggesting that all hadronic masses get their scale from the quark condensate, and therefore vanish at $T \to T_c$. However in a vacuum containing instanton

\textsuperscript{5}Note that the $N_f = 4$ case is missing, because the condensate is small and comparable to finite-size effects. Amusingly, recent results from Columbia group for $N_f = 4$ have found exactly this: a dramatic drop in chiral symmetry breaking effects (e.g. $\pi - \sigma, \rho - a_1, N - N^*(1/2^-)$ splittings.)
molecules it is not so, because these objects generate new non-perturbative inter-quark interactions. The Lagrangian describing those is discussed in [3]: in some channels like \( \pi, \rho \) it leads to attractive forces which (if strong enough) may create bound states even above the transition line. So far it is unclear whether it happens, but existence of these forces can be checked using the so called “screening masses”. Their T-dependence for a number of hadronic channels was calculated in the interacting instanton model, and those show good agreement with lattice ones. Especially important is strong attraction in scalar-pseudoscalar channels, shifting these masses down from their high-T asymptotic, \( M/\pi T = 2 \).

2. Looking for Equation of state in high energy heavy ion collisions

Recent data obtained with heavy ion beams (Au at Brookhaven and Pb at CERN) have displayed very strong collective flow effects (see other talks and my summary). These data can be used in order to obtain information about properties of the Equation of State (EOS) of hot/dense hadronic matter.

In order to do so, it is useful to return back from a (very complicated) cascade event generators to basic hydro description and can easily incorporate different scenarios (e.g., with or without the QCD phase transition). In contrast to longitudinal flow, for radial direction we know the initial conditions rather well, and therefore quantitative relation between the magnitude of the radial flow in central collisions and EOS can be made.

Let me show some results from (still unfinished) work by M.Hung and myself, which is aiming at developing a new model for AGS/SPS energy domain, called Hydro-Kinetic Model, HKM. The basic Equation of State (EOS) of hadronic matter used is that of a resonance gas, while for QGP one usually uses a simple bag-type EOS, with a constant fitted to \( T_c = 160 \text{MeV} \). In Fig.2(a) we show that phase boundary and the paths on the phase diagram. As both baryon number and entropy is conserved, the lines are marked by their ratio. Those for \( n_b/s = 0.02, 0.1 \) correspond approximately to SPS (160 GeV A) and AGS (11 GeV A) heavy ion collisions, respectively. Note that the trajectory has a non-trivial zigzag shape, with re-heating in the mixed phase. The endpoint of the QGP branch is known as the “softest point” [7], while the beginning of the hadronic one we will call the “hottest point” [6].

Experimentally observable particle composition is related to the stage of the collision known as a chemical freeze-out. Multiple works (e.g. [3]) have applied thermal description and determined where those points are on the \( T, \mu_b \) phase diagram: for both AGS and SPS those (inside error bars) coincide with the “hottest points” on our zigzag.

The next step is to define the effective EOS in the form \( p(\epsilon) \) (needed for hydro) on these lines: that is shown in Fig.2(b). Note that the QCD resonance gas in fact has a very simple EOS \( p/\epsilon \approx \text{const} \), while the “mixed phase” is very soft indeed. The contrast between “softness” of matter at dense stages and relative “stiffness” at the dilute ones is strongly enhanced for the SPS case.

\[ ^6 \text{Hydro does not contradict to event generators, but rather get support from them. It just allows to get space-time picture and flow much easier, without simulating all multiple re-scattering in more-or-less thermal conditions.} \]

\[ ^7 \text{Of course, in the “Hagedorn sense”, as the hottest point of the hadronic phase.} \]
The final velocity of the observed collective “flow” is time integral of the acceleration, which is proportional to $p/\epsilon$ ratio plotted in Fig.2(b). Although the observed velocity is not very different for both energies, this is kind of a coincidence since both the EOS and the time development of the collisions are rather different. In short, it was found that this basic EOS is too soft for AGS data on flow but describe well the SPS case. Most interesting, if one assume that there is no phase transition at all, and $p/\epsilon \approx .2$ like in the resonance gas, the flow obtained is way too strong.

The most difficult puzzle related to radial flow is provided by experimental data showing that it has very strong A-dependence (see e.g. talk by Gaardhoje in this proceedings): the larger the nuclei, the stronger is the flow. To resolve this (and other) puzzle one should correctly include the kinetics of the freeze-out. In most hydro papers, expansion was cut off at fixed $T$, usually about 140 MeV. In other words, it was assumed that the hadronic matter dries out very quickly, and its large pressure cannot lead to any significant motion. The situation is different if one apply correct kinetic condition: the (relevant) collision rate approach the expansion rate. Then one finds that with new high energy heavy ion beams (Pb at CERN and Au at BNL) we now have access to cooler hadronic gas, compared to medium ion collisions studied few years ago. The temperature of thermal freeze-out id only slightly smaller, 110-120 MeV, but in terms of space-time evolution this difference imply a very significant change. For sufficiently heavy ion collisions the matter gets the “extra push” from stiff pion gas at the end, which explains large flow velocity.

One interesting effect which follows from this observation: at the expansion stage without chemical but in thermal equilibrium the chemical potentials of all species change. The chemical potentials of the pions created in central PbPb it should reach $\mu_\pi = 60 – 80 MeV$. NA44 data indeed support extra low-$p_t$ enhancement in PbPb relative to S-induced reaction.

All phenomena discussed are predicted to be very much enhanced at RHIC. Larger multiplicity lead to (paradoxically) smaller final temperatures, and softness of the EOS due to the QCD phase transition change the lifetime of matter by a factor of 2.

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8It was also known from RQMD-based studies that extra repulsive interaction between baryons is indeed need to reproduce radial flow.

9It was originally pointed out by G.Baym, see details in 6.

10For clarity: those potentials are conjugated to total number of particles, so say for pions they enter distributions of $\pi^+, \pi^-, \pi^0$ with the same sign.
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Figure 2.  (a) Paths in the $T - \mu$ plane for different baryon admixture, for resonance gas plus the QGP; (b) the ratio of pressure to energy density $p/\epsilon$ versus $\epsilon$, for different baryon admixture.