Matter Formed at the BNL Relativistic Heavy Ion Collider

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We suggest that the “new form of matter” found just above \( T_c \) by RHIC is made up of tightly bound quark-antiquark pairs, essentially 32 chirally restored (more precisely, nearly massless) mesons of the quantum numbers of \( \pi, \sigma, \rho \) and \( a_1 \). Taking the results of lattice gauge simulations (LGS) for the color Coulomb potential from the work of the Bielefeld group and feeding this into a relativistic two-body code, after modifying the heavy-quark lattice results so as to include the velocity-velocity interaction, all ground-state eigenvalues of the 32 mesons go to zero at \( T_c \) just as they do from below \( T_c \) as predicted by the vector manifestation (VM in short) of hidden local symmetry. This could explain the rapid rise in entropy up to \( T_c \) found in LGS calculations. We argue that how the dynamics work can be understood from the behavior of the hard and soft glue.

Introduction— Within a month or two of operation, RHIC found the “new matter” they formed to be an extremely strongly interacting liquid, nothing like the quark gluon plasma.

This brings up two issues: a long-standing “old” issue as to what the structure of the state is in the vicinity of, and below, the presumed chiral phase transition point and what the proper tool to understand it is and a “new” issue as to what lies above the critical point to which the accepted theory of strong interactions, QCD, is supposed to be able to access perturbatively. In this Letter, we wish to address these issues in terms of an old idea on in-medium hadron properties proposed in 1991 [1] which has been recently modernized with the powerful notion of “vector manifestation (VM)" of hidden local symmetry theory [2] and buttressed with results from lattice QCD. Our principal thesis of this paper is that just as it required an intricate and subtle mechanism to reach the VM structure of chiral symmetry just below \( T_c \) – which is yet far from fully understood – from the standard linear sigma model picture applicable (and established) at \( T \sim 0 \), the structure of matter infinitesimally above \( T_c \) could also be vastly intricate and subtle from the starting point of QCD at \( T \sim \infty \) at which asymptotic freedom is applicable and “established.” We will make here an admittedly leaping conjecture by extrapolating and inferring from available information coming from lattice results and hinted by RHIC data that the structures of matter just below and just above \( T_c \) could be related. Further arguments to support this conjecture will be presented in a later publication.

In this paper we shall focus on a new state of matter formed at high temperature probed by RHIC experiments, particularly in the vicinity of the critical temperature \( T_c \sim 175 \text{ MeV} \). We should mention that the question of how matter changes from one form of symmetry realization to another form of symmetry realization is a general and fundamental issue in physics and in the case of hadrons, this still remains a more or less open question. Our conclusion drawn based on a variety of arguments to be developed below will be that the transition is “continuous” across \( T_c \) in the sense that the same degrees of freedom that come up from below \( T_c \) are found just above \( T_c \) although perhaps in a disguised form.

As we will develop in this paper, the key to the possible new matter produced at RHIC and its connection to the matter below \( T_c \) is in the glue from the gluons exchanged between quarks and its role in Brown-Rho scaling.

Two types of glue— The first observation we make is that there are two types of glue, a soft one and a hard one (epoxy). This aspect has been emphasized recently in [3]. The fact that there are two glues was first found in the thesis work at Columbia by Yuefan Deng [4]. He found in lattice gauge simulations (LGS) that about half of the total glue, the soft glue, was melted as \( T \) reached \( T_c \). The exchange of soft gluons holds hadrons together at temperatures below the phase transition temperature \( T_c = 175 \text{ MeV} \). Below this temperature their coupling to hadrons produces a soft gluon condensate which is melted as \( T \) goes up to \( T_c \). As the soft glue melts the constituent quarks turn into massless current quarks. The hard glue (epoxy) on the other hand begins to melt only above \( T_c \). It is the epoxy condensate that produces the length parameter \( \Lambda_{QCD} \) by breaking scale invariance and through what is called “dimensional transmutation.”

Brown-Rho scaling— In 1991 [1] Brown and Rho predicted using a dilaton field that as the soft glue melted with increasing temperature (or increasing density) hadron masses would decrease essentially proportionally to the decrease in the soft gluon condensate, going to zero at \( T_c \). Their scaling found in [1], \( \frac{m^*}{m} \sim \frac{f^*_q}{f_q} \), can be related to the quark condensate,

\[
m^* \sim \frac{\langle \bar{q}q \rangle^*}{\langle \bar{q}q \rangle}.
\]  

Two recent independent studies in a generalized hidden local symmetry approach, one by Harada and Sasaki [5] and the other by Hidaka et al. [6], confirmed that both the vector and axial vector mesons go massless at the chiral restoration satisfying [6].

This so-called "Nambu scaling" was confirmed in heat
bath by Koch and Brown \[7\] from lattice results. Thermodynamics of quasiparticles; \textit{i.e.}, that the entropy of a hadron with effective mass \(m^*\) can be obtained from that of mass \(m\), simply replacing \(m\) by \(m^*\), holds up. So entropy is just counting. The main correction to the early work is that the \(T_c = 140\,\text{MeV}\) used there should be increased to \(T_c = 175\,\text{MeV}\). The energy in the gluon condensate is therefore increased by \((140/175)^4 = 2.44\). This is important for the bag constant, which is just 1/4 of the gluon condensate.

There are experimental indications that vector meson masses do undergo Brown-Rho scaling in medium. It has been shown by Shuryak and Brown \[5\] that the results of the STAR detector at RHIC demonstrated that the \(\rho\) meson mass was somewhat decreased at low density at RHIC and a recent CBELSA/TAPS collaboration experiment \[6\] unambiguously showed that the mass of the \(\omega\) meson while inside a tin nucleus is also less than in free space. Both decreases follow the quantitative estimates of Brown and Rho with Nambu scaling, corresponding to an \(\sim 20\%\) drop in mass by nuclear matter density.

\textit{Meson masses are zero at} \(T_c\) --- It was argued by Brown et al. \[10\] that mesons whose masses decrease with temperature to zero at \(T_c\) could describe the large entropy increase found in lattice gauge calculations as \(T \rightarrow T_c\). (Massless mesons give a greater entropy than massive ones.) In other words, the phase transition was described as mesons going massless, that is, chirally restored. The spontaneous breaking of chiral symmetry which gives the mesons their scalar masses is restored at \(T_c\). It is plausible that the multiplet structure at \(T_c - \epsilon\) reflects Weinberg’s “mended symmetry” \[11\].

\textit{Thermal masses} --- One of the surprises of the Bielefeld lattice gauge calculations was that they found the quarks to have very large thermal masses, of the order of \(m_q \sim 1\,\text{GeV}\) at \(T = 1.5\,T_c\) and similarly large masses at nearby temperatures. These thermal masses arise from the self-energy diagrams at finite \(T\) and, in our scenario, are the only masses just above \(T_c\). It should be stressed that they are \textit{not} scalar masses which break chiral symmetry but energies, that is, fourth components of four vectors. It has not been possible to calculate these masses analytically since the interactions just above \(T_c\) are supposed to be strong and hence nonperturbative. Weldon \[12\] obtained in perturbation theory a somewhat complicated formula involving momentum for the dispersion relation which could however be approximated to within \(\sim 10\%\) by \(p_T^2 \approx m_q^2 + p^2\) with \(m_q\) a temperature dependent quantity which could be used as an effective mass in a simple hydrogenic model derived from Bethe-Salpeter equation as done in \[13\] following the argument of Hund and Pilkhun \[14\]. Remarkably, despite the presence of the effective mass, the helicity remains conserved as the quark wave function satisfies a free Dirac equation. We shall assume that this result holds non-perturbatively with the masses given by the Bielefeld lattice calculations.

Now \(T_c = 175\,\text{MeV}\) so that the Boltzmann factor for the quark \(e^{-m_q/T}\) is tiny. Similar results were found for gluons. Thus, there are not enough quarks and gluons present to produce the pressure, \textit{etc.} found in RHIC experiments. If the mesons \(\pi, \sigma, \rho, \alpha_1\) in \(SU(4)\) multiplets go massless at \(T_c\), they can furnish the pressure. But they can only go massless if the binding of \(\sim 1\,\text{GeV}\) quarks and antiquarks is \(\sim 2\,\text{GeV}\), so the bound states have to be small, \(\sim h/(2\,\text{GeV})\) in radius, much smaller than the typical \(h/(mc)\) of the usual mesons, where \(m\) is their mass. The calculation in \[13\] (described below) found the rms radius of the bound states just above \(T_c\) to be \(\lesssim 0.4\,\text{fm}\). The pion is an exception to this. The interaction must bring \(2m_q \sim 2\,\text{GeV}\) to zero, and therefore is very strong.

The Bielefeld group \[15\] have carried out lattice gauge calculations to obtain the heavy quark free energy and entropy for the region of temperatures above \(T_c\). Their results (for pure gauge) were used by Park, Lee and Brown \[16\] in agreement with their full QCD calculations \[17\] in almost all respects, once a rescaling by the relevant temperatures is made. From the energy and entropy the Coulomb potential \(V\) can be obtained

\begin{equation}
V(r,T) = F(r,T) - T \frac{\partial F(r,T)}{\partial T}.
\end{equation}

We shall pay particular attention to \(T \gtrsim T_c\) where we have \(V(r)\), suppressing the \(T\). The caveat here is that defining the thermal modification of a potential energy between the quark-antiquark pair is known to be complicated and subtle \[17\].

Now the calculations are for quarks of infinite mass. In order to use them for light quarks, which make up the \(SU(4)\) multiplets \(\pi, \sigma, \rho, \alpha_1\) we must put in the magnetic (velocity dependent) additions. We will be guided by Brown \[18\] who considered the velocity dependence of the interaction between the two \(K\)-electrons in heavy atoms. In the case of Uranium, \(Zu \sim 2/3\) so we are not so far from the \(\alpha_q \sim 1\) encountered here. Since the color singlet interaction is simply a Coulomb one, the difference from the Coulomb interaction in atomic physics being that since QCD is a gauge theory, it runs with scale, we can use the apparatus of high-Z atomic physics to put in the velocity dependence. A nice result of our procedure is that we obtain the correct number of degrees of freedom, that of 32 massless bosons as determined by LGS. The total interaction for stationary states is, translated to QCD,

\begin{equation}
V = - \frac{\alpha_s}{r} \left(1 - \vec{\alpha}_1 \cdot \vec{\alpha}_2\right),
\end{equation}

where \(\vec{\alpha}_1\) and \(\vec{\alpha}_2\) are the Dirac velocity operators. The \(\vec{\alpha}\)'s are also helicities, and have eigenvalues \(\pm 1\) above \(T_c\). Thus, since the quark and antiquark move in opposite directions, \(\vec{\alpha}_1 \cdot \vec{\alpha}_2 = -1\) and

\begin{equation}
V = - \frac{2\alpha_s}{r}.
\end{equation}
These results follow because the quark and antiquark are in helicity eigenstates \[\uparrow\] and \[\downarrow\].

Lattice gauge calculations above \( T_c \) show collective excitations S, P, V and A, all degenerate \[\uparrow\] and \[\downarrow\]. These are just the Nambu-Jona-Lasinio collective states (collective because they are a sum over quarks and quark-holes) which continue on up through \( T_c \) \[\uparrow\].

In \[\uparrow\] the masses of the bound states (all 32 are degenerate) above \( T_c \) were calculated with input of the lattice results for the Coulomb potential \[\uparrow\] \[\uparrow\] with magnetic effects added. Since, the (thermal) quark masses \( m_q \) have not been calculated at \( T_c \), (only at \( 1.5T_c \) and \( 3T_c \)) calculations were made for various assumptions within the range \( 1 < 2 \) GeV. It was found there that the meson masses go to zero at \( T_c \) regardless of this mass.

A simple hydrogenic model derived from Bethe-Salpeter equation \[\uparrow\] is invoked in \[\uparrow\] to show why the meson masses are zero at \( T_c \) irrespective of \( m_q \). The resulting equation is essentially the Klein-Gordon one since the spin-dependent interactions can be neglected because of the large inertial parameter in the denominator of the magnetic moment as discussed below. It may very well be that this parameter goes to \( \infty \) as \( T \) comes down to \( T_c \) from above as a result of confinement, making our approximation and the Weinberg mended symmetry \[\uparrow\] exact. In this model the heavy quark Coulomb potential is started from zero at \( r = 0 \) and increases to \( 2m_q \) (with string breaking ignored) at large distances. The chirally restored mesons would have energy \( \sim m_q \) in this approximation (half of the absolute value of the \( -2m_q \) potential energy). Introduction of the velocity–velocity interaction brings the meson masses to zero.

The chirally restored mesons have been found in quenched LGS at \( T \approx 1.4T_c \) and \( 1.9T_c \) \[\uparrow\] \[\uparrow\] \[\uparrow\] \[\uparrow\]. In the heavy quark approximation Park et al. \[\uparrow\] deduced that the binding energy was 0.15 GeV and the thermal mass \( m_q \sim 1.2 \) GeV at \( 1.4T_c \). Nonlinearities were not displayed; all 32 modes were degenerate. On the other hand, Brown at el. \[\uparrow\] found in the light quark formalism that just above \( T_c \) there were large nonlinearities. The width for \( \rho \to \pi \pi \) was estimated as 380 MeV. (Above \( T_c \) this is the only interspecies transition, the \( \pi \) and \( \sigma \) and the \( \rho \) and \( a_1 \) being equivalent). Possibly the nonlinearities produced enough noise in the lattice calculations to prevent them from being extended below \( 1.4T_c \). (Difficulties did appear \[\uparrow\].) In any case the large \( \Gamma(\rho \to \pi \pi) \) strongly interspecies interactions.

Rescaling at \( T_c \).— Just above \( T_c \), the color singlet (Coulomb) potential predominates completely over the colored interactions. Thus only the colorless states matter. This can be understood roughly in the following way.

As long as hadrons with masses are present, chiral symmetry is broken and scales are set by the chiral symmetry breaking one \( 4\pi f_\pi \approx 1 \) GeV, where \( f_\pi \approx 90 \) MeV is the pion decay constant. However, once chiral symmetry is restored at \( T_c \), the scale is determined by the \( \pi \), \( \sigma \), \( \rho \), \( a_1 \) meson (scalar) masses which are zero at \( T_c \). Suppose we describe this effect – which is clearly absent in perturbation theory and hence highly non-perturbative – in terms of an effective Coulomb interaction with an “effective charge” that simulates movement towards the infrared, so that the “Coulomb interaction,” evaluated in lattice gauge calculations of the Polyakov loop \[\uparrow\], becomes large, \( \alpha_s \sim 2 \). The effective coupling constant, with the \( 4/3 \) value of the Casimir operator and the factor 2 for velocity dependence included, comes to \( \alpha_s \sim 4/3 \) as \( T \to T_c \). Since \( \alpha_s = g^2/4\pi \) this means that \( g \to \sim 8 \), and \( g > 1 \) being the strong coupling. This strong coupling is manifested in the interactions for \( T \geq T_c \).

At first sight it may seem strange that the masses of the bound states which result from lattice calculations employing quite large bare (or current) quark masses \( \tilde{m} \) all drop so sharply toward zero. That explicit chiral symmetry breaking plays an insignificant role can be understood by realizing that in the presence of the explicit breaking, the total mass is

\[
M = \sqrt{m_{th}^2 + \tilde{m}^2} \approx m_{th} + \frac{\tilde{m}^2}{2m_{th}} \tag{5}
\]

with \( m_{th} \sim 1 \) GeV, so that \( M \sim m_{th} \) unless extremely high \( \tilde{m} \)’s are used in the LGS. Thus, the lattice calculations of the system entropy at \( T_c \) depends little on the explicit chiral symmetry breaking.

As the soft glue melts the interactions go to zero.— Let us now go up to \( T_c \) from below in the hadron sector. Since the hadronic interactions are given by the exchange of soft gluons, as these “melt” and disappear, the interactions go to zero. We have not yet rigorously established but we believe that the melting of the soft glue corresponds to the gauge coupling \( g_V \to 0 \), i.e., “the vector manifestation (VM)” in HLS theory \[\uparrow\]. In fact both the vector coupling constant \( g_V \) and the mass \( m_{th} \) go to zero at a fixed point at \( T_c - \epsilon \). This gives rise to “hadronic freedom.” This implies that as the fireball expands and \( T \) drops below \( T_c \) in heavy-ion processes, interactions between the hadrons cease, building up again as \( T \) drops to zero.

In going up to \( T_c \) from below, the dynamically generated masses of the mesons, such as the \( \rho \), go to zero leaving only the bare mass \( \tilde{m} \sim 5 \) MeV. Since the width of the \( \rho \) goes as \[\uparrow\], replacing \( m_{th}^* \) by \( 5 \) MeV one sees that the effects from the explicit chiral symmetry breaking are negligible. Thus, in conjunction with eq. \[\uparrow\], we see that the explicit chiral symmetry breaking changes the description of the chiral restoration transition only negligibly from that in the chiral limit; i.e. the effects from crossover, rather than second order phase transition would be difficult to see. We might have expected this from the fact that the behavior of the glue (see Fig.2 of \[\uparrow\]), which has no quarks in it, is a good guide to the chiral restoration transition.

We should note that whereas the VM with movement toward \( g_V = 0 \) at the fixed point brings the vector mesons to zero mass just below \( T_c \) and the symmetry determining
the zero mass of \( \pi \) and \( \sigma \) just above \( T_c \) is chiral symmetry (which protects the pion from mass), the apparent \( SU(4) \) symmetry which seems to set \( m_\pi^0 \) and \( m_\sigma^0 \) equal to \( m_\pi \) just above \( T_c \) is not exact. The magnetic moments of quarks and antiquarks above \( T_c \) are given by eq. (21) of Brown et al. [12]

\[
\mu_{q,\bar{q}} = \pm \sqrt{\frac{2\alpha_s}{\pi}} \frac{p_0}{m_{th}}
\]

where \( p_0 = E - V \), and \(-V \sim 2m_{th} \sim 2\text{GeV} \). The large inertial mass in the denominator of the spin dependence suppresses it greatly, but since the spin dependence is not zero, the \( SU(4) \) symmetry is not exact.

\textbf{Concluding Discussion.} — Although our principal focus was on the vicinity of \( T_c \), it is interesting to extend the scenario to higher \( T \). The four detector groups conducting research at RHIC have announced that they have created a new form of matter, “the perfect liquid”. The initial temperature at which RHIC matter is formed is \( \sim 2T_c \). This liquid follows from the scenario of Shuryak and Zahed [24], which can easily be understood from the results of [13]. Namely, the thermal masses, divided by \( T \), decrease toward the perturbative value

\[
\frac{m_{th}}{T} = \frac{g}{\sqrt{6}},
\]

so that the \( q\bar{q} \)-pairs will become unbound at \( T \sim 2T_c \). As they unbind their scattering amplitudes become very large in magnitude, going through \( \pm \infty \), so that the quark and antiquark mean free paths are small, giving the low viscosity of the perfect liquid.

Now there can be myriads of other states, both colored and colorless, which contribute substantially to the viscosity of the perfect liquid. The bound state masses in [13] using their results. GEB and BAG were partially supported by the U.S. Department of Energy under Grant No. DE-FG02-88ER40388 and part of the work of MR was supported under Brain Pool program of Korea Research Foundation through KOFST, grant No. 051S-1-9. .

\textbf{Note added} — After the paper was accepted for publication, it was brought to our attention that we had inadvertently left out an important paper that appeared in the same year as that of Ref. [4] by Su Hong Lee [25] who suggested the notion of two glues exploited in our paper and discussed also the possibility of nonperturbative quark-gluon dynamics being operative above \( T_c \).

References

[1] G. E. Brown and M. Rho, Phys. Rev. Lett. \textbf{66} (1991) 2720.
[2] M. Harada and K. Yamawaki, Phys. Rept. \textbf{381} (2003) 1.
[3] G. E. Brown, L. Grandchamp, C.-H. Lee and M. Rho, Phys. Rept. \textbf{391} (2004) 353.
[4] Y. Deng, Nucl. Phys. B (Proc. Suppl.) \textbf{9} (1989) 334. These calculations were renormalized, had the black-body radiation removed, etc., so that they could be put with QCD sum rule calculations by C. Adami, T. Hatsuda and I. Zahed, Phys. Rev. D\textbf{43} (1991) 921.
[5] M. Harada and C. Sasaki, \textsf{hep-ph/0511312}.
[6] Y. Hidaka, O. Morimatsu and M. Ohtani, \textsf{hep-ph/0512375}.
[7] V. Koch and G.E. Brown, Nucl. Phys. \textbf{A560} (1993) 345.
[8] E.V. Shuryak and G.E. Brown, Nucl. Phys. \textbf{A717} (2003) 322.
[9] D. Trnka \textit{et al.}, Phys.Rev.Lett. \textbf{94} (2005) 192303.
[10] G.E. Brown, A.P. Jackson, H.A. Bethe and P.M. Pizzochero, Nucl. Phys. \textbf{A560} (1993) 1035.
[11] S. Weinberg, Phys. Rev. Lett. \textbf{65} (1990) 1177; “Unbreaking symmetries,” in Trieste Proceedings “Salam-festschrift,” 1995.
[12] H.A. Weldon, Phys. Rev. D\textbf{26} (1982) 2789.
[13] H.J. Park, C.-H. Lee and G.E. Brown, Nucl. Phys. \textbf{A763} (2005) 197.
[14] V. Hund and H. Pilkuhn, J. Phys. B: At.Mol.Opt.Phys. \textbf{33} (2000) 1617.
[15] O. Kaczmarek, F. Karsch, F. Zantow and P. Petroczky, Phys. Rev. D\textbf{70} (2004) 074505; O. Kaczmarek, F. Karsch, P. Petroczky and F. Zantow, Phys. Lett. B\textbf{543} (2002) 41.
[16] O. Kaczmarek, S. Ejiri, F. Karsch, E. Laermann and F. Zantow, Prog. Theor. Phys. Suppl. \textbf{153} (2004) 287; O. Kaczmarek and F. Zantow, Phys. Rev. D\textbf{71} (2005)
[17] F. Karsch, Eur.Phys.J. C43 (2005) 35.
[18] G.E. Brown, Phil. Mag. 43 (1952) 467.
[19] M. Asakawa, T. Hatsuda and Y. Nakahara, Nucl. Phys. A715 (2003) 863.
[20] P. Petreczky, J. Phys. G30 (2004) S431.
[21] G.E. Brown, C.-H. Lee and M. Rho, Nucl. Phys. A747 (2005) 530.
[22] T. Hatsuda, private communication.
[23] G.E. Brown, C.-H. Lee, M. Rho and E. Shuryak, Nucl. Phys. A740 (2004) 171.
[24] E.V. Shuryak and I. Zahed, Phys. Rev. C70 (2004) 021901; Phys. Rev. D70 (2004) 054507.
[25] Su Houng Lee, Phys. Rev. D40 (1989) 2484.