The Ideal Liquid Discovered by RHIC, Infrared Slavery Above and Hadronic Freedom Below $T_c$

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Abstract

We construct the nature of the matter found in RHIC when its temperature has dropped down close to, and below, $T_c$. Just above $T_c$ it is composed of extremely strongly bound quark-antiquark pairs forming chirally restored mesons of the quantum numbers of the $\pi, S, \rho$ and $a_1$ with very small size and zero energy and just below $T_c$, it is composed of mesons of the same quantum numbers with zero mass. We invoke infrared slavery for the former and the vector manifestation (VM) of hidden local symmetry for the latter. As the temperature drops below $T_c$, the strongly bound quark-antiquark pairs are ejected into what is basically a region of “hadronic freedom” in which the interactions are zero. Experimental evidences for this are seen in the STAR data.
1 Introduction

The possible discovery of a new state of matter at RHIC [1] highlights a fascinating aspect of QCD in nonperturbative regime that has hitherto been unappreciated in the field. First of all, it is indicating that when hadronic matter makes the transition from Nambu-Goldstone mode to Wigner-Weyl mode at a critical temperature \( T_c \), it does not just go into the boring gas of liberated quarks and gluons as many in the field have anticipated or hoped but into something much more ordered and correlated that has many features in common with trapped atomic gases and dual black-hole horizons [2]. Perhaps even more significantly, it reinforces the notion that physics is continuous across the critical temperature as it has been noticed in the case of density with “quark-baryon continuity” [3]. How the continuity is effectuated is not yet understood. In this paper, we make the first step toward the unraveling of this intricate phenomenon and sketch how the RHIC discovery can be understood in terms of the in-medium scaling of masses and coupling constants developed since a long time as one goes up to \( T_c \) from below and of the strong-coupled gauge interactions recently uncovered as one comes down to \( T_c \) from above.

The starting point of our discussions is Fig. 1 obtained by Dave Miller [4] which contains a great deal of important information on how the phase change can take place. It shows clearly that two kinds of glue are involved in the chiral structure of hadronic vacuum, the soft and hard glues. The soft glue is locked to the quark condensate and responsible for the chiral phase transition and the hard glue responsible for the dimensional transmutation in QCD is left more or less unaffected by the phase change. This two-glue scenario is developed in a greater detail in the next section.

2 Development

We carry out our discussion within the framework of the Nambu-Jona-Lasinio (NJL) model as developed in [5]. One can think of the NJL as arising when the light-quark vector mesons (or more generally the tower of vector mesons as implied in dimensionally deconstructed QCD [6] or holographic dual QCD that arises from string theory [7]) and other heavy mesons are integrated out. As such, it presumably inherits all the symmetries of QCD and many of the results of hidden local symmetry theory that captures the essence of QCD [8]. In order to handle phase transitions in effective field theories, one has to treat properly the quadratic divergence which is present in loop graphs involving scalar fields. How this can be done in a chirally invariant way is explained in [8]. Now the cutoff that figures in the calculation represents the scale at which the effective theory breaks down. The natural scale for this is the chiral symmetry...
Fig. 1. Gluon condensates taken from Miller[4]. The solid line shows the trace anomaly for SU(3) in comparison with that of the light dynamical quarks denoted by the open circles and the heavier ones by filled circles.

breaking scale $\Lambda_{\chi_{SB}} = 4\pi f_\pi \sim 1 \text{ GeV}$. Brown and Rho [5] showed however that the cutoff in NJL should be at $4\pi f_\pi/\sqrt{2} \sim 700 \text{ MeV}$. This cutoff was
suggested to be suitable for Wilsonian matching to constituent quarks rather than to QCD proper, i.e., current quarks. For Wilsonian matching to QCD proper, the scale commensurate with vector mesons has to be incorporated into the theory. In addressing this problem, Harada and Yamawaki [8] were led to introduce hidden local gauge invariance which allowed the vector meson mass to be counted as of the same order in the chiral counting as the pion mass. In that case the loop of Fig. 2 comes in so as to cancel the $\sqrt{2}$ in the NJL cutoff denominator, so the Wilsonian matching radius is raised to $4\pi f_\pi$.

NJL was carried out most neatly by Bernard, Meissner and Zahed [9]. Their cutoff $\Lambda = 700$ MeV is close to $4\pi f_\pi/\sqrt{2}$. In BGLR [10] we got our best fits for $\Lambda = 660$ MeV and the NJL $G\Lambda^2 = 4.3$ which gave $T_c \simeq 170$ MeV. (Here $G$ is the dimensionful coupling constant.) In a mean-field type of mass generation, such as is shown in Fig. 3, it can be thought of as the coupling to constituent quarks of a scalar meson $S$ with free-space mass $\sim 700$ MeV, $G \sim -g_{sQQ}^2/m_s^2$.

At $n = 0$, $T = 0$, the proper variables are nucleons. They are bound states of three quarks, bound together by the glue. They have mass $m_N$. Scale invariance is broken by filling negative energy states with them down to momentum scale $\Lambda$. Thus

$^1$ We denote the would-be scalar chiral partner of the pion by $S$ and reserve $\sigma$ for the longitudinal component of the $\rho$ meson that figures in hidden local symmetry theory [8].
\begin{equation}
B(\text{glue}) = 4 \int_0^\Lambda \frac{d^3k}{(2\pi)^3} \left\{ \sqrt{k^2 + m_N^2} - |\vec{k}| \right\},
\end{equation}

where we have subtracted off the perturbative energy $|\vec{k}|$. The integral is easily carried out with the result

$$B(\text{glue}) = 0.012 \text{ GeV}^4,$$

the value usually quoted for QCD sum rules.

We note that there is no melting of the glue until $T = 125$ MeV. We can understand this by that the nucleon masses are just too heavy to be pulled out of the negative energy sea. But as the temperature $T$ is increased the nucleons dissociate into constituent quarks. For instance, Meyer, Schwenger and Pirner [11] use a wave function which can be written schematically as

$$\Psi = Z |N\rangle + (1 - Z^2)^{1/2}|3q\rangle.$$

This implies that there is a transition of nucleons dissociating into constituent quarks. At this stage the glue which surrounds the quarks starts to melt, and the curve for $G^2(T)$ drops rapidly, down to $G^2 \sim 0.0045$ at $T_c$(unquenched). The heavy filled squares are for bare quark masses which are 4 times greater than the open MILC-collaboration ones, but the glue – unlike hadron masses – is insensitive to explicit chiral symmetry breaking. We can estimate the amount of soft glue by changing variables from nucleons to constituent quarks, where our degeneracy factor is 12,

$$B(\text{soft glue}) = 12 \int_0^\Lambda \frac{d^3k}{(2\pi)^3} \left( \sqrt{k^2 + m_Q^2} - |\vec{k}| \right).$$

Taking $m_Q = 320$ MeV, we find $B(\text{soft}) \sim 0.5 (0.012) \text{ GeV}^4$. In the LGS the drop is a bit more than half of the $T = 0$ glue. This could be achieved by choosing a somewhat larger constituent quark mass, say, $m_Q \sim 400$ MeV.

Now at $T_c$(unquenched) we are left with $G^2(T) \sim 0.005 \text{ GeV}^4$. This is at $T_c$ where the soft glue has all melted and the constituent quarks have become (massless) current quarks. Note that the next point at $T \sim 1.4T_c$ is equally high. There is no melting of the glue between $T_c$(unquenched) and $1.4T_c$(unquenched). This is why we call this glue epoxy (hard glue). It makes up the (colorless) Coulomb interaction which strongly binds the quark-antiquark molecules above $T_c$. (Just above $T_c$ it binds them to zero mass.) The epoxy is the glue that one finds in pure gauge calculations.
A convenient picture of what might be going on here is offered by an instanton model. In the version discussed in [12], the soft glue is composed of random instantons, which flip quark helicity in scattering and consequently break chiral symmetry. As the temperature rises above $T_c$, the remaining instantons rearrange themselves into instanton molecules, held together by quark zero modes. The instanton molecules, which are the chirally restored mesons, are the normal modes of quark, antiquark and glue. See Fig. 4. Note that the $u$ will be Coulomb attracted to both the $d$ and the $\bar{u}$, and the $d$ to the $u$ and $\bar{d}$. Thus the Coulomb attraction acts as an additional clamp holding the instanton molecules together, which helps to explain why the glue acts as epoxy just above $T_c$, so tightly bound are the instanton and antiinstanton. As noted, the curve for $G^2(T)$ is completely flat just above $T_c$ (unquenched) = 175 MeV, indicating that the glue there is not being melted at all.

Of course, the hadrons below $T_c$ (unquenched) interact with each other by way of the soft glue exchange between quarks in the two interacting hadrons. As the soft glue melts, the interactions go to zero. It is reasonable that the interactions between hadrons go to zero at the same rate as the hadron masses. Indeed, this is confirmed in the vector manifestation of hidden local symmetry by Harada and Yamawaki [8]: They find that $m^*_\rho$ and $g^*_V$ go to zero at the same rate $\propto \langle \bar{q}q \rangle$ at the fixed point at $T_c$ as one comes up to $T_c$ from below.

The lattice calculations give us a somewhat more precise picture of the behavior far from the critical point than the RG flow of Harada and Yamawaki. Namely, there is no movement in the soft glue up to $T = 125$ MeV and then it is melted rapidly in the next 50 MeV to $T_c = 175$ MeV. We interpreted this as showing that the relevant variables were nucleons up to 125 MeV, at which temperature they loosened into constituent quarks. Since the latter were bound only by $m_Q^*$ plus kinetic energy in the negative energy sea, the melting of the constituent quarks began then.

On the whole, the general behavior of the soft glue in the LGS supports the Harada-Yamawaki picture [8]. However on the one hand, it is much more
detailed in the LGS. On the other hand, the general mathematical reason that it behaves as it does is not as clear. Therefore, the two descriptions, LGS and the vector manifestation of Harada and Yamawaki, complement each other. The vector manifestation which provides a firm theoretical support to Brown-Rho scaling [13] predicts that the masses of the the 32 $\bar{q}q$ bound states go to zero as $T$ goes up to $T_c$ from below. From the LGS results of the Bielefeld group, Park, Lee and Brown [14] have shown that the masses of the same 32 $\bar{q}q$ bound states, which are just chirally restored mesons, go to zero as $T$ goes down to $T_c$ from above. There is a sort of continuity in physics between $T_c - \epsilon$ and $T_c + \epsilon$ for infinitesimal $\epsilon$. What governs this continuity is not understood at all.

It is commonly thought that the explicit chiral symmetry breaking changes the phase transition into a crossover one, substantially changing the behavior. We argue that this is not the case. In BLR[15] we showed that the width of the $\rho$-decay into two pions went as $m_\rho^5$. Thus, at chiral restoration, where the $\rho$ mass is only of order the bare mass $\bar{m}_\rho$ (that is, the mass stripped off of chiral condensates), the width becomes

$$\frac{\Gamma^*(T_c)}{\Gamma_\rho} \sim \left(\frac{\bar{m}_\rho}{m_\rho}\right)^5$$

coming up to $T_c$ from below. This is a completely negligible crossover.

It is remarkable that the scenario sketched above was captured in the 1993 paper by Brown, Jackson, Bethe and Pizzochero[16]. Both LGS and the HLS/VM are providing a compelling support to this scenario. They argued that mesons rather than liberated quarks and gluons are the correct variables up to well beyond $T_c$(pure glue) = $T_c$(quenched) and that above $T_c$(unquenched)= $T_{\chi SB}$ the (hard) glue remained condensed. Thus, although mesons could be regarded as quark-antiquark pairs, each quark must be connected with an an antiquark by a “string” (i.e., a line integral of the vector potential) in order to preserve gauge invariance. It was then noted that it is difficult to include the important consequences of such quark/antiquark correlations in the thermodynamics of the gluon condensed systems unless one continues to regard mesons as the correct effective degrees of freedom even above $T_{\chi SB}$. The fact that the string couples quark and antiquark so tightly, as a Coulomb potential, that “infrared slavery,” defined below, resulted as we understand now greatly simplified the Brown et al. picture [16].

In Kaczmarek et al. [17] the authors note “that at sufficiently short distances the free energy agrees well with the zero temperature heavy quark potential and thus also leads to a temperature independent running coupling. The range of this short distance regime is temperature dependent and reduces from $r = 0.5$ fm at $T = T_c$ to $r = 0.03$ fm at $T = 12T_c$.” Since the rms radius of
our chirally restored molecule just above $T_c$ is only 0.2 fm it is clear that the molecule is bound by the Coulomb force but if it were not, it would be bound by the confining force, as operates for $T < T_c$. Thus the LGS show that the same confinement below $T_c$ persists well above $T_c$, the chiral restoration having no effect until distances become much larger than those of the molecules just above $T_c$. In other words, the strings between quark and antiquark actually run the show, controlling the thermodynamics, for a sizable range of temperatures above $T_c$. It is no wonder, then, that as $T$ goes below $T_c$ with introduction of the order parameter $4\pi f_\pi \sim 1$ GeV which enters with the breaking of chiral invariance that the infrared slavery above $T_c$ will change to hadronic freedom below. Whereas Brown et al. [16] did not anticipate the hadronic freedom just below $T_c$, Brown et al. [18] ran into trouble with the Hagedorn energy which gets lowered by dropping masses to $T_{Hagedorn} \sim 0.75 T_c$. This would give two scales to the problem, which is inconsistent. This difficulty is satisfactorily removed if the hadronic interactions go to zero as $T \to T_c$.

### 3 Going above $T_c$: Infrared slavery

In Eq. (5) we found that going up to $T_c$ from below, the hadronic interactions went to zero at $T_c$ in the chiral limit, as the soft glue melted or equivalently approaching the vector manifestation fixed point, and that even with explicit symmetry breaking through the current quark mass, these interactions were still completely negligible. Thus what we might call “hadronic freedom” – more precisely defined in the next section – is reached just below $T_c$.

Now the pion is protected from gaining mass by chiral invariance and the $S$ is degenerate with the $\pi$ at $T_c$, so we can go up through $T_c$ in such a way that the $\pi$ and $S$ masses remain zero. This is not so easy as Hatsuda and Kunihiro thought [19] because of the strong coupling gauge interaction that we call “infrared slavery.” In other words, with the change in order parameter $4\pi f_\pi \sim 1$ GeV with chiral restoration to zero, the masses of the $\pi$ and $S$ remain zero. (And, essentially those of $\rho$ and $a_1$ because their masses, although not quite zero just above $T_c$, are very small due to the large thermal masses. These enter into the denominator of the magnetic moment operator (Eq. (21) of Brown et al.[20]); i.e.

\[
\mu_{q,\bar{q}} = \pm \sqrt{\alpha_s/p_0} 
\]

where $p_0 = E - V$ and $-V \sim 2m_{\text{thermal}} \sim 2$ GeV.) We return to the large thermal mass below. Practically speaking, one has the degeneracy of $\pi, S, \rho$ and $a_1$, 32 degrees of freedom, just above $T_c \, ^2$, although the $\rho$ and $a_1$ are not

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2 It would be more satisfying if an underlying symmetry argument were found to
Fig. 5. The large distance behavior of $\alpha_s(T)$ from evolution of the Polyakov loop in quenched LGS\cite{17}. Casimir factor ‘$4/3$’ and color magnetic interaction factor ‘$2$’ are not included.

united by any fundamental symmetry with the $\pi$ and $S$; they are important simply because they lie very close to zero energy. These 32 degrees of freedom precisely match the number of massless boson modes necessary for the entropy calculated in LGS, thus resolving the long-standing puzzle.

The most significant development in our understanding of the state just above $T_c$ is the large thermal mass and the way the thermal mass is balanced by the strong Coulomb attraction. Why are the thermal masses so large? We propose that the thermal mass is large because there is “infrared slavery.” Put in physical terms, the color Coulomb force becomes very large and the coupling of gluons to the quarks and of gluons to gluons in the cloud about a quark become very large. This can be seen in Fig. 5 which shows the large distance behavior of $\alpha_s(T)$ from the Polyakov loops in quenched QCD \cite{17}. In this figure $\alpha(T) = g^2/4\pi \rightarrow 2$ as $T \rightarrow T_c$ from above. This looks large already but, indeed, this isn’t even the half of it: BLRS \cite{20} have shown that the correction of heavy quark Coulomb interactions up to light quark ones brings the interaction to

$$V = -\frac{\alpha_s}{r} (1 - \vec{\alpha}_1 \cdot \vec{\alpha}_2) \quad (7)$$

be operative to make all 32 degrees of freedom degenerate and massless at $\epsilon$ above $T_c$ but we do not have any idea what that can be.
where the $\alpha$’s are the Dirac matrices. Including the Ampère’s Law interaction in $-\alpha_1 \vec{\alpha}_1 \cdot \vec{\alpha}_2 / r$ doubles the attractive Coulomb interaction for the quarks and antiquarks of opposite helicity and sets it to zero for those of the same helicity. The net result is that out of the initial 64 quark-antiquark degrees of freedom, 32 of them are brought down substantially in energy. Including the Casimir factor of $4/3$, this doubling brings

$$\alpha_s(T_c) = \frac{16}{3},$$

or $g = 8$, which is indeed a very strong coupling. This is how infrared slavery manifests itself. In fact, Park, Lee, and Brown [14](PLB) have shown, using the Bielefeld LGS essentially for full QCD, that the color Coulomb potential, when put into a relativistic two-body equation, brings the masses of the chirally restored mesons $\pi, S, \rho, a_1$ (just the quark-antiquark bound states) essentially to zero energy at $T_c$.

4 Going below $T_c$: Hadronic freedom

These colorless chirally restored hadrons then regain their full freedom as the temperature drops below $T_c$; in fact, they propagate as free noninteracting particles because their interactions have gone to zero as $T \to T_c$ from below in accordance with the HLS/VM [8]. We attributed loosely “hadronic freedom” to this phenomenon. This notion can be sharpened by looking at the renormalization group structure of Harada-Yamawaki HLS theory. In this theory, ultraviolet completion is made by matching the HLS effective theory to QCD at a suitable matching scale $\Lambda_M$ near the chiral scale. The full theory so defined is quantized below $\Lambda_M$ by RGE and loop corrections. Now the RG flow for this theory consists of coupled equations for the hidden gauge coupling constant $g$, the parameter $a = f^2_\sigma / f^2_\rho$ (where $f_\sigma$ is the decay constant of the Goldstone boson $\sigma$ that corresponds via Higgsing to the longitudinal component of the massive $\rho$) and the pion decay constant $f_\pi$. Harada and Yamawaki [8] found that the gauge coupling runs with the one-loop beta function that goes as $-\frac{N_f}{2(4\pi)^2} \frac{87 - \alpha^2}{6} g^4$ (where $N_f$ is the number of flavors). The parameter $a$ also runs, so the flow structure of the theory is more complicated than that of QCD but $a$ remains less than 87 so the beta function is always negative. This means that $g$ has an “ultraviolet fixed point” $g = 0$, that is, the theory is in a peculiar sense asymptotically free. It turns out that $a$ goes to 1 at the same fixed point. This point coincides with the vector manifestation fixed point to which the system is driven as the critical point ($T$ or $n_c$) is approached from below. As $g$ goes to zero, hadronic interactions go to zero and hence the term “hadronic freedom.” It is worthwhile pointing out that
this is a feature unique in strong-coupling many-body systems, hitherto not seen in particle physics.

Note that this scenario of the chirally restored mesons going massless as \( T \to T_c \) from below was already employed by Brown et al. \[16\]. It was noted in \[15\] that there is evidence of the complete equilibration of the \( \rho \)-mesons at \( T = T_c + \epsilon \), then a long period in which they are essentially noninteracting until the temperature has fallen to nearly that for thermal freezeout in the STAR data.

We have, of course, dealt only with the collective modes of mesons, \( S, \pi, \rho, a_1 \); These are just the mesons calculated in Nambu-Jona-Lasinio below \( T_c \).

5 Hydrogenic orbits at \( T_c \)

In the heavy quark approximation; i.e., with the thermal quark and antiquark masses \( m_{th} \), PLB \[14\] showed that the binding energy of the orbits of all 32 degrees of freedom were \( m_{th} \) as \( T \to T_c \) from above. In other words, the potential energy was \(-2m_{th}\) and the kinetic energy was \( m_{th} \), just as in the hydrogen atom. The small spin dependence that separated the \( \rho \) and \( a_1 \) was not seen in their results. However, the doubling of the Coulomb interaction just doubled the binding energy, sending the masses to zero at \( T_c \). This was done in the chiral limit, but explicit chiral symmetry breaking should have little effect. This is because the breaking is carried out by a world-scalar mass \( \tilde{m} \) which is small compared with the thermal mass \( m_{th} \). Since the latter is the 4th component of a 4-vector, the total mass is

\[
m = \sqrt{m_{th}^2 + \tilde{m}^2} \approx m_{th} + \frac{\tilde{m}^2}{2 m_{th}} \tag{9}
\]

and the term from explicit symmetry breaking is down by \( \tilde{m}/m_{th} \).

We have a simple argument for why the masses come out to be zero at \( T_c \). In drawing the curve of potential vs distance between quark and antiquark we should begin from zero at \( r = 0 \) because of asymptotic freedom. (We have checked that our phenomenological regularization at short distances discussed in BLRS \[20\] is essentially the same as the lattice regularization at short distances.) Ignoring string breaking, our potential must increase up to \( 2m_{th} \), the sum of the energies of quark and antiquark. From our above arguments, this would then give a binding energy of \( m_{th} \) with the heavy quark potential in the hydrogenic approximation. With the factor of 2 in potential obtained in going to the light quark approximation, the binding energy is also doubled.
and the masses are brought to zero. This simple picture describes the results of the full calculation.

6 Conclusion

We note that in the entire range of temperatures covered by RHIC, up to the initial temperature $\sim 2T_c = T_{\text{zero binding}}$ where the mesons become unbound, the material found in RHIC is composed of mesons, as already argued in 1993 by Brown et al. [16]. In the hadronic sector these are the 32 mesons in the spin-isospin SU(4), although the multiplet energies are badly broken. These 32 mesons are evolved as collective states (bubble sums of positive energy quarks and quark holes in the negative energy sea) in the Nambu-Jona-Lasinio formalism. Now the same 32 mesons go through $T_c$. The $\pi$ and $S$, both with zero mass, go smoothly through in mass, as determined by chiral symmetry. The $\rho$ and $a_1$ have zero mass coming to $T_c$ from below, because of the VM, alias Brown-Rho scaling. They have a very small mass just above $T_c$, because their magnetic moments have the inertial parameter $p_0 \sim 2m_{th}$ in their denominators, and the thermal mass $m_{th} \gtrsim 1$ GeV is very heavy. In the lattice calculations higher in energy, say $T = 1.4T_c$, there appears to be an SU(4) symmetry to the vibrations representing the mesons [21], but our analysis shows the $\rho$ and $a_1$ to be (imperceptibly) above the $\pi$ and $S$ in energy.

Thus, the 32 degrees of freedom have zero mass coming up to $T_c$ from below (Brown-Rho scaling) and, aside from the imperceptible shifts of $\rho$ and $a_1$ mentioned above, they are all massless going down to $T_c$ from above. It was shown that the effects of explicit chiral symmetry breaking are negligible. As pointed out above, these 32 degrees of freedom have precisely the entropy found in LGS, to within the (high) accuracy of the latter.

Going up in temperature from $T_c$ the masses of the $\pi, S, \rho$ and $a_1$ grow in unison, and their binding energy decreases, until the mesons are liquidated by
breaking up into quark and antiquark at $T_{\text{zero binding}} \sim 2T_c$, in the region of initial RHIC temperatures. The liquidation of the mesons results in very strong interaction between the quarks and antiquarks, the scattering amplitudes going from $\infty$ to $-\infty$ during the breakup arriving at a conformal invariance fixed point. These strong interactions result in the perfect liquid recently subscribed to by the four collaborations (Brahms, Phenix, Phobos and STAR) [22].

We, on the other hand, believe that RHIC, considering now the phenomena created with increasing temperature, has taken our 32 mesons through a fascinating panorama. First, with increasing temperature they have lost their mass (Brown-Rho scaling) and lost their interactions (vector manifestation) going into hadronic freedom as $T$ goes up to $T_c$ from below. Then the selfsame mesons resurface just above $T_c$ in infrared slavery. Only at the region of the initial (highest) temperature reached by RHIC are they liquidated. Even that cannot be reconstructed because their density is so high and their orbits are so intertwined that they may be just bumping into each other as tightly packed neighbors [15]. Thus we agree that it’s a liquid and Shuryak and Zahed [23] have made the case for that, but only at the death of our mesons.

This brings us to the following anecdote: A lady once asked the famous philosopher Immanual Kant “Does everything end with death ?” Kant replied, “No, then begins the litigation.” So now we have had the litigation. But it’s not as interesting as the life before.

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Dedication

Two of the authors (G.E.B. and C.-H. Lee) would like to dedicate this article to their deceased collaborator Hans Bethe. “And at the age of 83, he apprenticed

\[ i.e., \text{in 1993}; \text{see Ref. [16]}. \]
himself to Gerald E. Brown of the State University of New York, Stony Brook, in order to learn lattice gauge theory, which predicts how nuclear matter is transformed at extremely high temperatures into a plasma of particles called quarks and gluons, is one of the most challenging in all of physics.” Journal of Chemistry and Engineering, March 21, 2005, p.57. “I’m interested in learning new things”, Bethe explained.

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