**Quantum Chromodynamics
and
Nuclear Physics at Extreme Energy Density**

Progress Report

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1. QUARK-GLUON PLASMA AND HIGH-ENERGY COLLISIONS

1.1. Thermal Photons and Dileptons from Chemically Nonequilibrated Quark-Gluon Plasma

(M. T. Strickland)

The subject of this work is the thermal emission of photons and dileptons from a non-equilibrium quark-gluon plasma. In particular I have studied the possibility of using these signals to determine plasma initial conditions under the model proposed by Biro, Muller, et al. My conclusion from this work is that these signals can be used to determine the initial conditions to within 0.25 GeV if these signals are separable from the background processes only to 3 GeV. If, however, the backgrounds can be subtracted to 5 GeV then accurate (0.05 to 0.1 GeV) determination of the initial conditions is possible. The figure below is a plot of the fractional yield of thermal dileptons and photons as a function of M and p_T for different stages of plasma evolution. Each consecutive line from bottom to top represents approximately 0.6 fm/c elapsed. From this figure we see that dilepton and photon production for large p_T and M is dominated by early time scales. Therefore, this region provides a record of the earliest stages of plasma evolution. On the other hand, for low energies production is spread evenly over all time scales thus concealing information about the initial state. It is for this reason that background subtraction becomes crucial.

Fig.1. (a) Percent yield of dileptons as a function of invariant mass for different stages of thermal evolution. (b) Percent yield of photons. Both are at RHIC energies, but similar results hold at LHC energies.

My future research will address the importance of medium-effects on processes that take place in a quark-gluon plasma. I plan to study medium effects on associated (or jet) production of photons and dileptons. Associated production is the production of particles due to jet fragmentation, and a cursory analysis shows that this production could be

suppressed by an order of magnitude at LHC energies due to in-medium effects. This is an important question since it has been shown that associated production of dileptons and photons is the dominant background process for large p_T and M at LHC energies. If we can extend the background subtractable region of the thermal yields we consequently increase our ability to determine initial conditions.

1.2. Random Branching Processes and Multiparticle Production

(S. G. Matinyan [1])

The general theory of random branching processes is involved for establishing the relation between the parameters k and $\bar{n}$ of the negative binomial distribution of the produced particle multiplicity:

$$P_n^{(k)} = \frac{\Gamma(n+k)}{\Gamma(n+1)\Gamma(k)} \frac{\left(\frac{\bar{n}}{k}\right)^n}{\left(1 + \frac{\bar{n}}{k}\right)^{n+k}}$$

for the produced particle multiplicity in high-energy hadron interactions. The relation

$$k = \frac{a\bar{n}}{\bar{n} - b},$$

where $a \approx 3$ and $b \approx 7$ integration constants was found from comparison with experimental data. It provides a satisfactory description of the data on multiplicity distributions in $pp(\bar{p}\bar{p})$ -collisions energies up to 1.8 TeV, and can be used to make predictions for higher energies.

These predictions are:

- i) restoration of a KNO scaling at high energy,
- ii) all higher moments C_q ($q = 3, 4, \dots$) of the multiplicity distribution are rising and saturate asymptotically,
- iii) the scaled peaks of the distributions are moving to the left toward its asymptotic value

$$z_{\text{peak}} \equiv \frac{n_{\text{peak}}}{\bar{n}} = 0.67 + \frac{2.27}{\bar{n}}.$$

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1.3. Pair Production in the Quantum Boltzmann Equation

(J. Rau)

The evolution of the quark-gluon plasma, believed to be formed in the course of relativistic heavy-ion collisions, is commonly described by means of a transport equation [1]. It is well understood how a transport equation can account for acceleration in external

fields, scattering, or (hadro-) chemical reactions of the microscopic constituents. There is, however, another physical process which becomes increasingly important at high energies: regions of very large chromoelectric field strength may develop and subsequently decay by emitting quark-antiquark pairs [2]. This gives rise to the fragmentation of chromoelectric flux tubes ('strings'), a mechanism frequently invoked to model hadron production [3]. To account for such spontaneous creation of particles, the transport equation has to be modified by a source term.

This source term has been derived [4] for the case of e^+e^- -creation in a constant homogeneous electric field, and evaluated numerically [5]. It is convenient to introduce $a := \hbar q E / \epsilon_\perp^2$ and to consider, rather than the source term itself, the dimensionless quantity $\eta(a, p_\parallel / \epsilon_\perp) := (\epsilon_\perp / q E) \exp(\pi / 2a) \dot{n}_{\text{eou}}(\vec{p})$. The distribution of the longitudinal momenta of the produced electrons is shown in figure 1, for various values of a . Clearly it is *not* narrowly peaked around $p_\parallel = 0$. Rather, electrons are being produced predominantly in the direction of the external field ($p_\parallel > 0$). Electrons moving in the opposite direction ($p_\parallel < 0$) are being annihilated: for them, the production rate is negative. As another surprising feature, η displays (approximately) periodic oscillations whose period scales with a . This may be understood qualitatively if one views pair creation as a tunnelling process from the negative to the positive energy continuum. The barrier between these continua has a spatial width of the order $\epsilon_\perp / q E$, inducing a 'momentum quantisation'

Fig. 1. The re-scaled production rate η as a function of $p_\parallel / \epsilon_\perp$ for weak fields ($a = 0.2, 0.3, 0.7$).

$\Delta p_\parallel \sim \hbar q E / \epsilon_\perp$ and thus $\Delta(p_\parallel / \epsilon_\perp) \sim a$. Interference of multiply reflected electron wave functions then leads to the observed oscillations.

The source term is found to be generally non-Markovian. It exhibits a characteristic

memory time

$$\tau_{\text{mem}} \sim \frac{\hbar}{\epsilon_{\perp}} + \frac{\epsilon_{\perp}}{qE}$$

which, for strong fields ($a > 1$), may lead to temporary violations of the H -theorem.

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1.4. QCD Evolution Equations for Partons in Nuclear Matter

(K. Kinder-Geiger and B. Müller)

Recently, considerable progress has been made in a better understanding of the space-time structure of parton interactions during the early stage of these reactions. The conclusion emerging from different independent investigations, is that—for RHIC energies and beyond—most of the entropy and transverse energy is presumably produced already during very early times by frequent, mostly inelastic semihard parton collisions involving typical momentum transfers of only a few GeV. The underlying notion is that the early stage of nuclear collisions at sufficiently high energies can well be described in terms of the space-time evolution of many interrelated parton cascades [1] based on renormalization group improved perturbative QCD and relativistic kinetic theory.

When the parton density becomes so large that the quanta begin to overlap in phase-space, recombination (fusion) processes become relevant and the density must saturate towards its limiting value. An important question is therefore: how is “QCD in medium” modified as compared to “QCD in vacuum”? In trying to gain a more quantitative knowledge about the microscopic parton dynamics in medium, the most urgent questions concern (i) the initial conditions regarding the parton substructure of large nuclei, in particular the small x region and the magnitude of nuclear shadowing effects; (ii) the role of color screening and color diffusion; (iii) the impact of the Landau-Pomeranchuk-Migdal effect; (iv) the space-time dependence of parton interactions with regard to the influence of the characteristic interaction times of parton scatterings and the formation times for gluons emitted in bremsstrahlung processes.

Recently it was pointed out by McLerran and Venugopalan [2] that a consistent perturbative calculation of parton structure functions at small values of the Bjorken variable x becomes possible when one considers the limit of a very thick nuclear target. A projectile propagating at high energy through such a target “sees” a very large area density of valence quarks, and hence experiences an effective screening of color interactions in the transverse direction. The condition for the applicability of perturbative QCD then is that the screening distance is much shorter than the confinement scale Λ^{-1} . This is satisfied for a sufficiently thick target at sufficiently small x .

We make use of this insight in order to explore the evolution of a parton cascade inside nuclear matter under conditions where perturbative QCD applies because medium induced effects, such as color screening and rescattering, provide dynamical cut-offs on a scale short compared to Λ^{-1} . In principle, our approach applies to the propagation of fast partons in any kind of dense medium, be it a thermalized quark-gluon plasma or ground state nuclear matter. The equations derived by us can therefore be applied to jet quenching in a QCD plasma as well as to the fragmentation cascade of a quark after deep-inelastic scattering in an infinitely large nucleus. For simplicity we have started by considering the following idealized problem. Beginning with a prescribed initial distribution of fast partons injected into infinitely extended nuclear matter by some highly localized process of space-time extent, $(Q_0^2)^{-1/2} \ll \Lambda^{-1}$, we want to follow the evolution of the parton distribution in laboratory time, as it propagates through the medium.

We attempt to establish a connection between the semiclassical probabilistic picture of parton evolution in the leading logarithmic approximation (LLA) and the time development of parton cascades in six-dimensional phase space within the framework of non-equilibrium kinetic theory. To do so, we need to clarify two fundamental issues: First we need to relate the Altarelli-Parisi-Lipatov (APL) evolution equations, which determine the change of the parton number densities under variation of the variables rapidity y and transverse momentum $k_\perp$, or $x \approx \exp(y)$ and $Q^2 \approx k_\perp^2$, with the Boltzmann equation, which controls the time evolution of the phase space densities in both momentum and coordinate space. Second, we must relate the experimentally accessible number densities of quarks, antiquarks, and gluons with the single-particle phase space densities.

The first point—the connection between the APL evolution equations and the Boltzmann equation—has been previously suggested [3] in the formal context of hadron structure which is rather different from our considerations of parton cascades in nuclear matter. Nevertheless, this showed that the APL equations can be derived in a statistical manner similar to the Boltzmann equation, by taking into account both the gain *and* the loss of partons due to successive $1 \rightarrow 2$ branchings in (x, Q^2) -space. This self-contained detailed balance eliminates the necessity of calculating vertex and wave function renormalization explicitly, because the loss terms naturally take over this role. As a consequence, the resulting evolution equations are free of divergences and satisfy the constraints imposed by momentum and quark number conservation automatically.

The measured parton number densities $a_N(x, Q^2)$, where $a \equiv q, \bar{q}, g$, give the probability for finding a quark, antiquark, or gluon, inside a nucleon with fraction $x = k_z/P$ of the longitudinal nucleon momentum P and with virtuality Q^2 , or transverse momentum $k_\perp^2$. Here $Q^2 \approx k_\perp^2$ sets the scale of hardness that is identified with the momentum transfer of an interaction of the parton with a weakly interacting probe (e.g. a virtual photon) that measures the nucleon substructure. At present, the parton number densities are experimentally accessible only in a space-time integrated way and therefore must be interpreted as instantaneous distributions of partons inside a nucleon. On the other hand, in statistical many-particle systems the phase-space distribution $F_a(E, \vec{k}; \vec{r}, t)$ is the probability density for finding a parton of species a in a phase-space element $d^3k d^3r$ at time t . Evidently F_a contains explicit additional information about the space-time structure of the initial

state nucleons or nuclei, which is only present in an averaged way in the measured parton number densities.

We extend the probabilistic approach to the space-time evolution of partons in nuclear matter. To do so, we first relate the Q^2 -evolution to the development with time t , and then we include not only the $1 \rightarrow 2$ branching processes, but also the reverse $2 \rightarrow 1$ fusion processes and in addition $2 \rightarrow 2$ scattering processes. The latter two types of processes indirectly also give rise to additional stimulated branchings. Stimulated emission, fusion and scattering processes are naturally absent in vacuum, but in medium they are indispensable ingredients for obtaining a complete set of transition amplitudes and a self-consistent parton evolution. In order to keep track of the off-shell character of partons between interactions, we introduce an additional parameter, the *age* τ of a parton, which is inversely related to its virtuality Q . Interactions, such as scatterings, can rejuvenate a parton but promoting it back to a higher degree of virtuality. Without interactions, the parton "ages" by getting closer and closer to mass shell, without ever reaching it since gluons never propagate as free particles.

The full evolution equation for the off-shell gluon distribution, including free propagation, branching, two-gluon fusion, and gluon-gluon scattering terms, reads [4]:

$$\begin{aligned}
\frac{\partial}{\partial t} g(x, \tau, t) = & -\frac{\partial}{\partial \tau} g(x, \tau, t) \\
& -\frac{1}{2} \int_0^1 dz \gamma_{g \rightarrow gg}(z) \frac{\overline{\alpha}_s(\tau)}{2\pi\tau} g(x, \tau, t) \\
& + \int_0^1 \frac{dz}{z} \gamma_{g \rightarrow gg}(z) \frac{\overline{\alpha}_s(\tau)}{2\pi\tau} g\left(\frac{x}{z}, \tau, t\right) \\
& - \int_{-1}^0 dx_2 \hat{g}(|x_2|) \int dp_{\perp}^2 \frac{d\hat{\sigma}_{gg \rightarrow gg}}{dp_{\perp}^2} \rho_N g(x, \tau, t) \\
& \times \frac{|x|P}{\tau^2} \frac{|x|P}{\tau^2} \int_{-1}^0 dx_2 \hat{g}(|x_2|) \frac{d\hat{\sigma}_{gg \rightarrow gg}}{dp_{\perp}^2} \Big|_{p_{\perp}^2 = |x|P/\tau} \\
& \times \rho_N \int_0^\infty d\tau_1 g(x_1, \tau_1, t) \left[1 + \theta\left(\frac{\tau_1}{\tau} - \frac{x_1}{|x|}\right) \right] \\
& + \int_{-1}^0 dx_2 \hat{g}(|x_2|) \int^{x_1 P/\tau} dp_{\perp}^2 \frac{d\hat{\sigma}_{gg \rightarrow gg}}{dp_{\perp}^2} \rho_N g(x_1, \tau, t) \\
& - \int_{-1}^0 dx_2 \hat{g}(|x_2|) g(x, \tau, t) \rho_N \Gamma_{gg \rightarrow g}(x, x_2, x+x_2) \frac{4\pi\alpha_s(M^2)}{M^2} \Big|_{M^2 = |x+x_2|P/\tau} \\
& + \int_{-|x|}^0 dx_2 \hat{g}(|x_2|) g(x-x_2, \tau, t) \rho_N \Gamma_{gg \rightarrow g}(x-x_2, x_2, x) \frac{4\pi\alpha_s(M^2)}{M^2} \Big|_{M^2 = |x|P/\tau}
\end{aligned}$$

The modified equation no longer is a differential equation in the variable $(\ln t)$, but rather in t directly, because the presence of interactions with the medium defines a characteristic time scale $(\sigma\rho)^{-1}$, the mean free time between scatterings, which breaks the scale invariance

of the fragmentation cascade. Attempts to solve the new evolution equations for gluons and quarks are underway.

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1.5. Pion Fusion in Peripheral pp Collisions

(A. J. Schramm, B. Müller)

One opportunity afforded by a Full Acceptance Detector at hadron colliders [1] is the ability to measure total and elastic meson-meson cross sections, using leading particle tags to isolate one-meson exchange contributions. In a pp or $p\bar{p}$ collider, virtual pions associated with the beams can fuse; this process can be isolated by requiring two forward-going quasi-elastically scattered Δ baryons that can be identified by their decay products. In peripheral pp collisions, this would lead to the process shown in figure 1,

$$pp \rightarrow \Delta + \Delta + \pi + \pi \rightarrow \Delta + \Delta + X$$

A possible theoretical approach to such events uses an "equivalent pion approximation", developed by analogy with the Weizsäcker-Williams method of virtual quanta in electrodynamics. The approximation is based on the idea that the meson field accompanying an energetic baryon consists of virtual quanta close to mass shell, and as such can be replaced with an equivalent pulse of real pions if the invariant mass of the two pion system is much larger than $4m_\pi^2$.

We applied these same ideas to the fusion of virtual pions, using techniques developed in some of our earlier work [2]. For proton beams A and B with center of mass energy s , the cross section for $\pi\pi$ fusion can then be written in the form

$$\sigma_{AB}^{\pi\pi} = \int dx_1 dx_2 f_\pi^A(x_1) f_\pi^B(x_2) \sigma_{\pi\pi}(s_{\pi\pi}),$$

where x_i is the fractional beam momentum carried by the virtual pions and $\sigma_{\pi\pi}$ is the production cross section in the collision of two real pions with squared center of mass energy $s_{\pi\pi} \equiv x_1 x_2 s$. In the language of the parton model, $f_\pi(x)$ is the quasi-elastic pion distribution function of the proton. In order to calculate f_π , the correct factors for the pion-baryon vertices are needed. Since the pions are taken to be close to mass shell, we used the results of heavy baryon chiral perturbation theory for low energy pion-baryon interactions.

Baryon chiral perturbation theory was developed by Jenkins and Manohar [3] using a variation of the formalism developed for the study of the heavy quark limit of Quantum Chromodynamics [4,5]. This approach assigns to each meson-baryon vertex a factor of

$$(c_{\pi BT}) \frac{C}{\sqrt{2}F_\pi} k^\mu ,$$

where $F_\pi = 93$ MeV is the pion decay constant, k^μ is the pion momentum, and $c_{\pi BT}$ is the Clebsch-Gordan coefficient at the πBT vertex. The constant $C = 1.53$ has been determined [3] by fitting the experimentally measured $T \rightarrow B\pi$ decay modes. In addition to this vertex factor, there will be a contribution from the spin sum over the spinor and Rarita-Schwinger fields. For the interaction shown in figure 1 the invariant matrix element is

$$\mathcal{M} = \bar{T}_\mu(p'_1) k_1^\mu B(p_1) A_{\pi\pi \rightarrow X}^I \bar{T}_\nu(p'_2) k_2^\nu B(p_2) ,$$

where $A_{\pi\pi \rightarrow X}^I$ is the amplitude for two pions of total isospin I to fuse into anything.

In our work we used these methods to study the process $pp \rightarrow \Delta^{++} + \Delta^{++} + X$, requiring the fusion of two π^- mesons; this reaction can be triggered on the four positively charged decay products of the Δ^{++} baryons. We calculated the total cross section and determined its dependence on the impact parameter of the pp system, taking nuclear absorption into account within the framework of the Glauber approximation.

Our results indicate a cross section of approximately $11 \mu\text{b}$; including absorption effects reduces this $\pi\pi$ cross section to about $5.25 \mu\text{b}$. This is roughly 0.025% of the total elastic pion-pion cross section. A luminosity of $10^{31} \text{s}^{-1} \text{cm}^{-2}$ would yield about 50 events per second; an integrated luminosity of 10pb^{-1} produces some 50 million events.

The majority of the nuclear absorption is expected to occur at small values of the impact parameter b . Indeed, the quasi-elastic nature of the interaction is maintained only in those collisions in which the two protons pass by each other. Thus it is important to verify that a significant portion of the $\pi\pi$ cross section extends out to relatively large impacts parameters. The dependence of the differential cross section on impact parameter is shown in figure 2, where the dotted curve shows the suppressing effects of absorption.

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2. HADRON STRUCTURE AND CHIRAL DYNAMICS

2.1. Hadron Structure Studies

(R. P. Springer)

We use the decuplet of baryons as a laboratory for probing hadronic structure. This is a rich field not only because the Continuous Electron Beam Accelerator Facility (CEBAF) will be collecting data in the near future on the Delta particle and other resonances in the baryon decuplet [1], but also because the octet baryons play a dramatic role in properties of the decuplet and vice-versa, which provides vital information on each. Because the hadronic scale exists in an energy region where the strong coupling constant is large, perturbative calculations in that coupling are not useful and, in lieu of an exact solution to Quantum Chromodynamics (QCD), one must find some small quantity to be an expansion parameter. This is provided by the joint concepts of chiral symmetry and the heavy baryon limit of hadron physics [2]. These together form a theory for low momentum processes ($E \leq 500\text{MeV}$) involving the lowest lying multiplets of mesons and baryons. It is model-independent in the sense that it depends only on the symmetries of QCD and takes any needed parameters from experimental measurements.

In our last Progress Report, we discussed the one loop corrections to strong decays of the decuplet, which seem to require that the coupling constants respect a spin-flavor symmetry not present in the QCD Lagrangian; the radiative decays of the decuplet, where in particular the SU(3) violating rates can be used to determine the importance of meson loops in these processes; and the electric quadrupole to magnetic dipole ratios of decay moments, which probes the structure of the decuplet members. We have extended these studies as described below, where the notation and operators are taken from the previous report.

2.2. Extracting $E2/M1$ Moments for the Δ .

(R. P. Springer)

One difficulty with the ratio of electric quadrupole to magnetic dipole moments, $\delta = \sqrt{\frac{5}{3}} \frac{A_{E2}}{A_{M1}}$, is that the interference of resonant with nonresonant terms makes it ambiguous to measure and to define. For the Δ , pion-photoproduction data is used, along with model-dependent extractions, to give an estimate of this [4]. An alternative approach is provided by the angular correlation between production and emission photons available through polarized Δ decay. CEBAF naturally produces different numbers of particles in the $m = \pm 3/2$ than in the $m = \pm 1/2$ state. Our result is

$$S_{\gamma\gamma}(\theta) = 1 + \left(\frac{1 + 6\text{Re}(\delta) - 3|\delta|^2}{2(1 + 3|\delta|^2)} \right)^2 P_2(\cos \theta) \quad (1).$$

where θ is the angle between the incoming and outgoing photons and $P_2(x) = \frac{1}{2}(3x^2 - 1)$ is the second order Legendre polynomial. $S_{\gamma\gamma}(\theta)$ is shown in Fig. 1. A measurement of $S_{\gamma\gamma}(\frac{\pi}{2})/S_{\gamma\gamma}(0)$ should give a determination of δ , with an estimate of the nonresonant pieces found by measuring the ratio at a variety of energies. This will complement the pion-photoproduction determinations of the δ ratio for the Δ particle.

Fig.1. A polar plot of $S_{\gamma\gamma}$ for $Re(\delta) = \delta = 0, 0.05, 0.1$ and 0.2 .

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2.3. Static Moments of the Decuplet and the Role of Strangeness.

Static moments can be used to unravel the structure of nonfundamental objects. In our case, we apply this technique to the first resonance of baryons to see if chiral perturbation theory will provide a deeper understanding of the hadrons. To one loop order in meson exchange, we have predictions for magnetic and quadrupole moments of the decuplet.

A tree level operator, in addition to a loop diagram, contributes to the magnetic moment of the decuplet. We find the coefficient of this operator by using the measured magnetic moment of the Ω^- hyperon. The magnetic moments of the neutral baryons do not depend on this counter term and are therefore the most sensitive to purely loop effects. Because we include the strange mesons as possible loop particles, a comparison of our predictions with measured moments will indicate the importance of strange loops in even non strange baryons (such as the Δ). Our results are shown in Fig. 1. The magnetic

moments satisfy contracted $SU(N_f)$ relationships, which are demanded by the large N_c limit of QCD. This points out the compatibility, at least for these processes, of the chiral and large N_c limits of QCD. (N_f is the number of flavors, N_c the number of colors.)

Fig.1. The magnetic moments of the decuplet baryons, in units of nuclear magnetons. The widths are associated with the uncertainties in the experimentally extracted coupling constants.

The lowest order quadrupole moments occur, in chiral perturbation theory, as one loop diagrams; no counterterm is necessary to this order, so that predictions are sensitive purely to meson loop effects. The pion loop graphs generally dominate over kaon loops because both generate a logarithm of the meson mass over the chiral symmetry breaking scale (~ 1 GeV). Yet, Clebsch-Gordon coefficients conspire to make the contributions from pion loops exactly cancel for isospin zero objects. This suggests that the first contribution from decuplet intermediate states for these objects will involve kaon loops. Therefore the deuteron, for instance, will not get a contribution to its quadrupole moment from any

Fig.2. The quadrupole moments of the decuplet baryons, from one-loop graphs. The dark points are the moments derived from the central values of the coupling constants and the lighter lines are the associated uncertainties.

decuplet baryons (the Δ is often considered in this regard) with pion loops, to this order. The dominant component of the quadrupole moment of the decuplet baryons is proportional to the third component of isospin. These moments are shown in Fig. 2. We expect these predictions to be tested at CEBAF in the coming years, and for measurements of the static moments to help test the validity of the heavy baryon chiral perturbation theory approach to understanding low energy QCD.

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2.4. Disorientated Chiral Condensates through Annealing

(B. Müller)

The order of the chiral restoration phase transition in QCD for realistic values of the up, down and strange quark masses is currently unknown. If the transition is nearly second order, Rajagopal and Wilczek [1] have speculated that the nonequilibrium dynamics of a heavy ion collision can generate metastable domains in which the chiral condensate is misaligned. Such domains can provide a coherent source of pions that can exhibit novel "centauro-like" fluctuations of neutral and charged pions [2]. Clearly, the ability of the coming ion-ion experiments at the RHIC and LHC colliders to resolve such fluctuations critically depends on the size and energy content of these domains.

Numerical simulations [3] show that the zero-temperature linear sigma model advocated in [1] leads to domains that are roughly pion sized, too small to be resolved in experiments. Nevertheless, we illustrate in this Letter that finite temperature effects in a heavy ion collision can enhance the size of domains relative to the zero-temperature estimates.

Rajagopal and Wilczek propose a "quench" scenario in which the condensate is initially chirally symmetric as appropriate at high temperature, but its evolution is taken to follow classical equations of motion in the absence of a heat bath. For this quench scenario to be realistic, the expansion rate of the heat bath needed to create the symmetric initial state must greatly exceed the rate at which the mean field evolves. The opposite is more likely the case. In [1] and [3], the σ and π fields in the $O(4)$ linear sigma model are taken to characterize the dynamics of the quark antiquark condensate in two-flavor QCD. These fields evolve over the time scale $\sim m_\sigma^{-1} \sim 1/3$ fm, where $m_\sigma = 600$ MeV is the zero-temperature σ mass. On the other hand, model calculations indicate that chiral symmetry is broken very late in RHIC-energy collisions, perhaps at proper times $\tau_c \sim 5 - 20$ fm, so that the expansion time scale $\sim \tau_c$ is much larger in comparison.

We discuss an alternative "annealing" scenario in which the sigma-model condensate evolves semiclassically in the presence of a bath of quasiparticle excitations [4]. Fluctuations induced by the quasiparticles create an effective potential in which a chirally symmetric state is initially stable (modulo finite quark-mass effects). As the system expands and

rarefies, the effective potential develops a “wine-bottle” shape and slowly changes towards its free-space depth. The condensate then evolves in this changing potential. Transient domains of disoriented $\Phi \equiv (\sigma, \vec{\pi})$ condensate can develop because the $\Phi \approx 0$ initial state is unstable in the wine bottle potential. The system “rolls down” from the unstable local maximum of $V_0(\Phi)$ towards the nearly stable values with $|\Phi| = v$, since the symmetry breaking term $-H\sigma$ is relatively small. This process is analogous to spinodal decomposition in condensed matter physics. Field configurations with $\vec{\pi} \neq 0$ develop during the roll-down period. The field will eventually settle into stable oscillations about the unique vacuum $(f_\pi, \vec{0})$ for $H \neq 0$, but oscillations will continue until interactions eventually damp the motion. In the heavy-ion system a domain can radiate pions preferentially according to its isospin content.

Domain growth in the annealing scenario with an evolving effective potential differs from the quench as follows. Initially, V_{eff} is nearly flat for $\sigma \approx \vec{\pi} \approx 0$. Therefore, the roll-down time scale can be very large at first, or infinite if the corresponding equilibrium transition is truly second order (as is certainly the case for $H = 0$). Only as the potential approaches its free space shape does the roll-down become rapid. In both quenching and annealing, domains can grow as long as σ is substantially different from f_π . In the annealing scenario, however, that time scale is limited by the time needed for V_{eff} to approach V_0 . That time in turn depends on the global dynamics of the nuclear collision.

We describe the evolution of the condensate in the context of a self consistent Hartree-like semiclassical approach. There, the condensate evolves as a classical field in the presence of a fluctuating quasiparticle bath. We neglect the effect of collisions among quasiparticles, and the dynamical effect of the quasiparticle masses. To illustrate how the fields and quasiparticles evolve with time, we restrict our attention for the moment to a single domain in the interior of the collision volume. An expanding system created in a central Au+Au collision at RHIC will reach temperatures below $T_c \approx 132$ MeV only after a time $\sim 5-20$ fm comparable to the Au transverse radius $R_A \sim 7$ fm. We take the quasiparticle flow to be roughly 3-dimensional and homologous.

The numerical results [4] depend on the strength of the Hartree term at the critical temperature, expressed by the parameter ϵ . The value $\epsilon = 1$ applies if chiral restoration in an equilibrium system occurs as a true second order phase transition. Domains can grow as long as the system is unstable, i.e. until the σ field begins to oscillate about the value $\sigma \sim f_\pi$. For a nearly critical system with $\epsilon = 1$, we estimate this time to be $\tau_f - \tau_c \approx 1.5 \tau_c$, see figs. 2. Such values imply that the domain size can reach $R_D \approx 7$ fm — very large indeed! We obtain a more conservative estimate $R_D \sim 3-4$ fm by taking $\epsilon = 0.5$ and supposing that domain growth stops after σ reaches f_π for the first time at $\tau_f - \tau_c \approx (0.3 - 0.5)\tau_c$. Note that the numerical effect of the inflation factor $(\tau_c/\tau)^2$ is insignificant.

We have shown that annealing can lead to domains significantly larger than those expected from a quench. Our results therefore strengthen the suggestion that disoriented chiral condensates can be observed in heavy ion collisions at RHIC.

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2.5. Time-Dependent Variational Calculation in Disorientated Chiral Condensates

(Chiyong Lin and B. Müller)

In high multiplicity hadronic collisions it was suggested that the observed pions can be produced by the coherent decay of the classical pion field. Some authors have speculated that this classical pion field is oriented along a single direction in the isospin space disoriented from its external value, so that the condensate is associated with the isovector rather than the isoscalar degree of freedom. The signature for this event is the large number of either neutral pions or charged pions. This phenomenon may have been observed in the Centauro event in cosmic ray collisions. Rajagopal and Wiecek have proposed the nonequilibrium dynamics for the chiral transition to explain such phenomenon. They use a sigma model to describe the collective behavior where its time evolution follows the classical equations of motion at zero temperature. Here we propose to investigate this problem using time-dependent variational principle with gaussian trial wave functional as ansatz. This technique has been shown to be a powerful tool to explore the off-equilibrium situation in quantum field theory, such as the dynamics of the scalar field in the inflationary scenario. We use this approximation to linear sigma model and obtain a set of self-consistent equations of motion for the pion and sigma field. Analytical and numerical investigation are currently in progress.

2.6. Pion Condensation in Relativistic Heavy Ion Collisions

(C. Greiner [*], C. Gong and B. Müller)

Rajagopal and Wilczek [1] recently speculated that in a rapid quench through the second order chiral QCD phase transition disordered chiral configurations may emerge. The scenario explicitly assumes large deviations from thermal equilibrium. The spontaneous growth and subsequent decay of these configurations would give rise to large collective fluctuations in the number of produced neutral pions compared to charged pions. These kind of fluctuations can also be understood as follows [2]: For very energetic collisions the central rapidity region is presumably free of baryons and thus contains no isospin quantum number. If now a major fraction of the produced pions settles in the same state, the wavefunction for an overall vanishing total isospin can be explicitly constructed. The probability of finding n neutral pions in such a coherent state falls like $\sim 1/\sqrt{N_{\pi^0}}$, so that

the highest probability occurs for events with a low number of neutral pions. In addition, we also found that particular interesting correlations among the various pion pair channels then should occur which also could be seen in future experiments at RHIC.

Regardless of this speculative picture of decaying chiral condensates, it is useful to ask whether a condensed pion state can be formed under circumstances where the hadronic matter never gets far out of thermal equilibrium or if the phase transition occurs not spontaneously, but slowly, as one would may expect because of the surplus in latent heat. If a locally thermalized state is created throughout the rapidity gap between projectile and target nuclei, can the density become so high that some fraction of the pions are forced into a condensate? The measured slopes in the pion spectra at the CERN-SPS show a

Fig. 1. The entropy per pion s_π/n_π is plotted as function of temperature T for several fixed chemical potentials μ_π . Note, that the ratio s/n decreases for smaller temperatures at the critical chemical potential $\mu_c(T)$ (dashed line).

remarkable enhancement at low transverse momentum, which, in fact, may be explained by introducing an effective chemical potential of $\mu_\pi \approx 90 - 130$ MeV near the free mass of a pion, indicating that during the reaction the pions evolve out of chemical equilibrium. Condensation would occur if μ_π reaches the pion mass. The driving out of chemical equilibrium is motivated by the fact that the annihilation rate of pions becomes small at still rather high temperatures, so that in the ongoing cooling and expansion the number of (effective) pions remains more or less constant. In addition, the pions 'feel' an attractive interaction in the ρ -channel, which, as our calculations show, effectively lower the pion mass in the hot hadronic environment. On a first sight one may expect that this would present an opportunity for the onset of condensation.

In this (ideal) picture the entropy to pion ratio should thus be constant during the expansion of the fireball. It now turns out, however, that this ratio along the line following the critical chemical potential *decreases* with falling temperature (see Fig. 1). Hence, because the total entropy can never decrease, the pion system can not condense in a forthcoming evolution, if the hadronic matter stays in thermal equilibrium throughout the cooling process. Only if the hadronization out of a QGP is spontaneous (on the time scale of inelastic collisions), a fraction of pions could settle in a Bose condensated state and this state may survive until decoupling.

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2.7. Quark Pair Production in the Chiral Phase Transition

(C. Greiner [*])

One particular goal in relativistic heavy ion collisions is the formation of a deconfined quark gluon plasma. It is expected that this phase transition is accompanied by a restoration of chiral symmetry. The massive (confined) quasiparticles, the constituent quarks, become bare, undressed quarks with current quark masses much below the constituent quark masses. However, from the onset the reaction during such a collision proceeds far off equilibrium, a thermal QGP may be reached only after some finite time in the further evolution process. The 'undressing' of the quarks should occur in between this evolution. Being still far from equilibrium this restoration of chiral symmetry may happen rather spontaneously. The scalar quark condensate $\sigma \sim \langle \bar{q}q \rangle$ would melt on a time scale smaller than any scale set by the dynamics. Denoting the time of occurrence at $t = 0$, the light quark masses then are a function of time

$$m(t) = \begin{cases} m_c \approx 350\text{MeV} & \text{for } t \leq 0 \\ m_b \approx 10\text{MeV} & \text{for } t > 0 \end{cases}$$

Such a sudden change of the underlying vacuum structure would result in the production of quark and antiquark pairs. The constituent quark vacuum becomes unstable so that during the restoration of the masses occupied states of the negative Dirac continuum are partially transferred with some specific amplitude into the open positive continuum. The amplitude for this is obtained in a straight forward way within the in/out formalism. A particle with momentum $\mathbf{p}$ sitting in the negative Dirac sea is described by the incoming wavefunction $\psi^{(\text{in})} = \varphi_{\mathbf{p}\downarrow}^{(c)}(\mathbf{x}, t)$. After the spontaneous change of the vacuum structure this wavefunction goes over into $\psi^{(\text{in})} \rightarrow \alpha_{\mathbf{p}} \varphi_{\mathbf{p}\downarrow}^{(b)} - \beta_{\mathbf{p}} \varphi_{\mathbf{p}\uparrow}^{(b)}$, where the states with the label b denotes the outgoing (bare) wavefunctions. The quark (pair) production rate is then obtained as the projection squared on the appropriate outgoing states defining the particles [1] and takes here the simple form

$$N_{\mathbf{p}} = \langle 0, \text{in} | \hat{N}_{\mathbf{p}}^{(\text{out})} | 0, \text{in} \rangle = |\beta_{\mathbf{p}}|^2.$$

In this simplified model of a spontaneous vacuum decay this probability can be calculated analytically by solving the one-particle Dirac equation [2]. It is found that the overall quark production resembles the distribution of a saturated hot quark plasma with a temperature between 200 and 300 MeV, although the distribution does not drop off exponentially. For bare quark masses in the range of 150 MeV, like for the strange quarks, the production is diminished for low momenta.

We thus conclude that if for dynamical reasons a transition from the broken to the unbroken chiral phase occurs rapidly, the change of the underlying vacuum structure is accompanied by a nonperturbative production of quarks and antiquarks. It would be interesting to know how the numbers are changed if the transition occurs more smoothly in some finite time interval [2].

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2.8. Gluon Assisted Pair Production in the Chiral Phase Transition (C. Greiner [*])

In this work we want to continue our examination of the consequences of a spontaneous change of the vacuum when restoring chiral symmetry. According to the results of the parton cascade model the quarks will reach a chemical saturation only eventually whereas the gluons saturate on a significantly smaller timescale (on the order of half a fm/c). The reason for this behaviour lies in the fact that the quarks are only produced in second order by gluon annihilation and the rate for this process is small compared to the corresponding one for gluon production. The decay of one single gluon into a quark and antiquark is, of course, kinematically not allowed, as long as all partons are treated as effective on-shell particles. During the above mentioned transition, however, the quarks are no more on-shell particles, their spectral function contains a wide off-shell spectrum. This gives raise to the idea that a single gluon could in principal decay into a quark-antiquark pair during the decay of the vacuum. In return a large fraction of energetic gluons thus may (spontaneously) decay into such pairs and would contribute to an accelerated chemical saturation of the quarks in phase space.

The problem can be rephrased as follows: What is the number of quarks (and antiquarks) produced in a first order decay of already present gluons sitting in a time dependent scalar background field? The formulation we have carried out by using real-time Green functions [1] being appropriate for non-equilibrium studies. The number of produced quarks is contained in the '<'-component of the complete one-particle Green function and is given by the projection on the corresponding outgoing state in the far future

$$N_p = \lim_{t_1=t_2 \rightarrow \infty} \int d^3x_1 d^3x_2 \left((-i) \varphi_{p \uparrow}^{(b) \dagger}(1) G^<(1,2) (\gamma_0 \varphi_{p \uparrow}^{(b)}(2)) \right) .$$

For our purpose we have to specify now the Green function. The perturbative interaction with the gluons are diagrammatically straightforward. The first order decay is contained in the self-energy insertion Σ for the quarks. Casting the contributions from the time contour path to real times, the first order correction reads

$$\Delta G^<(1,2) = \int d^4x_3 d^4x_4 \left\{ \begin{array}{l} G_0^{\text{ret}}(1,3) \Sigma^{\text{ret}}(3,4) G_0^<(4,2) \\ + G_0^{\text{ret}}(1,3) \Sigma^<(3,4) G_0^{\text{adv}}(4,2) \\ + G_0^<(1,3) \Sigma^{\text{adv}}(3,4) G_0^{\text{adv}}(4,2) \end{array} \right\}.$$

where the 0-th order Green function G_0 is determined nonpertubatively by the full set of the incoming states involving in time according to the time dependent scalar background Hamiltonian. Specifying the individual self-energy terms the number of produced particles can now be calculated [2]. Two reasons, however, makes the analysis very cumbersome. The first being purely technical lies in the fact that the Green functions depend on both time arguments separately due to the time dependend external field. The usual Feynman techniques can hence be not applied. The second stems from the occurrence of IR singularities. Besides the pair production the above expression also incorporates the absorption or emission ('Bremsstrahlung') of gluons for the quarks being produced already from the (0-th order) vacuum decay. A separation of the pair production and these processes is not possible because of quantum mechanical interference terms. In the end one has to show that these singularities cancel, as they do. In addition to these two difficulties there also arises an UV infinite contribution out of the (pertubative) gluonic vacuum. Its interpretation lies in the usual mass renormalization. (Also there exists a finite mass renormalization contribution due to the interaction with the already present real gluons). However, being interested only in the contributions due to the presence of real gluons, the vacuum terms can explicitly be separated and will thus be discarded.

For a numerical study we parametrize the incoming gluons by a thermal distribution. This may not be too physical because the gluons may still be distributed far from equilibrium; yet it is convenient to see how large the final numbers are and by interpreting these to reason what might then be expected for other distributions. The study shows that a large fraction of the highly energetic gluons will indeed decay into quarks and that the slope of the produced quarks falls off exponentially, dictated more or less by a temperature the same as used for parametrizing the gluon distribution [2]. Such a phenomena could give raise to a nonpertubative dynamical creation of quarks and antiquarks (as well as entropy) in ultrarelativistic heavy ion collisions. More numerical studies are presently underway.

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2.9. Multiply Strange Nuclear Systems

(C. Greiner [*])

The formation of so called strangelets, small pieces of strange quark matter (SQM), had been proposed as a 'unique' signature for the formation of the QGP [1] in present and future heavy ion collision experiments. Because of their expected large strangeness content they would occur as rather neutral and unusual heavy 'hadronic' particles in the detectors.

Until recently the assumption that strangelets could be the only metastable configurations with such an unusual small charge to mass ratio seemed to be very plausible. However, one could in principle also think of small strange nuclei, which are made out of nucleons and strange baryons. If one thinks for example of a state $(2p, 2n, 2\Lambda, 2\Xi^0, 2\Xi^-)$, such a state would possess a baryon number of 10, it would be neutral and contains the same amount of strange quarks as either up and down quarks. (In fact, according to our new understanding, such a multistrange nuclei turns out to be bounded and thus metastable.) In addition, because all of the various baryons are located in the $1s$ -state, its average density of such a configuration is also comparable to that inside a strangelet, being almost 3 times that of normal nuclear matter. The detection of a nearly neutral massive 'particle' would then not be an unambiguous signal for the discovery of a strangelet. However, on the other hand, the possibility of having new forms of nuclei may turn out to be as exciting.

Essential for the stability of such new nuclear systems is the understanding of the Y-N and Y-Y interaction inside those multihypernuclei. The starting point is the description of the single hypernuclei, where one neutron is substituted by a strange Λ . From the measured one-particle states of the hyperons the individual potential depths and the strength of the attractive Y-N forces are extracted quite accurately. The experimental access to the Y-Y forces is, however, more delicate. There exist so far only three candidates for double-hypernuclei. Yet a careful analysis has given that the residual interaction of the two hyperons, more specifically the $1s$ -matrix element $\Delta B_{\Lambda\Lambda} = \langle 1s | V_{\Lambda\Lambda} | 1s \rangle \approx -5$ MeV, is also attractive and seems to be nearly as strong as the N-N-interaction inside nuclei. (A phenomenological consideration, which postulates a general SU(6) symmetry among the various potentials, supports this reasoning - see Fig. 1.) The residual interaction can now be simulated by introducing two new mesons which only couple to the hyperons. An explicit relativistic shell model calculation reveals that indeed the experimental data on the double hypernuclei can be successfully reproduced [2]. Moreover, a whole 'zoo' of exotic multihypernuclear systems is now being postulated.

Such strange hadronic matter (SHM) shows remarkably quite similar global properties as SQM. The only quantitative difference might lie in the overall binding energies of such exotic configurations. If the binding energy lies, however, in the same region, it turns out that presumably slightly positive candidates would be rather multihyperonic states and that slightly negative candidates may be both states [3].

Fig.1. Neutron-neutron (nn) and hyperon-hyperon potentials $V(r)$ for the 1S_0 state in a one-boson-exchange model.

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3. NONPERTURBATIVE STUDIES OF NONABELIAN GAUGE THEORIES

3.1. Lyapunov Spectrum of Classical SU(2) Lattice Gauge Theory (C. Gong)

A rescaling method has been developed for SU(2) lattice gauge theory and allows following the evolution of the system for a long time as well as obtain other Lyapunov exponents besides the largest one. But due to practical reasons, the complete Lyapunov spectrum for the theory can only be obtained on a lattice of small sizes [1]. Results are shown in Figure 1.

Fig. 1. Lyapunov spectrum in SU(2) gauge theory on lattices with different sizes N . The solid line corresponds to a 3^3 lattice; the dashed line is for a 2^3 lattice. Only the positive Lyapunov exponents are shown. The exponents λ_i are scaled with the maximal Lyapunov exponent $\lambda_{\max} = \lambda_1$ for each lattice size, and the index i is scaled with total number of Lyapunov number $18N^3$.

In the figure, The solid line shows the Lyapunov spectrum on a 3^3 lattice; the dashed line is for a 2^3 lattice; The points correspond to the result on a 1^3 lattice. The spectra are plotted in the scaled form: the Lyapunov exponents are scaled with respect to the

largest Lyapunov exponents; the indices are scaled with respect to the total number of Lyapunov exponents, which is $18N^3$ on a N^3 lattice. The spectrum is nicely separated into three parts: one-third of positive exponents, one-third zeros and one-third negative ones. The general feature of anti-mirror symmetry of a Lyapunov spectrum of a Hamiltonian system is observed. The number of zeros equals the number of Gauss' law conservations on the lattice, which implies, at least on the lattice, there is no hidden symmetry besides the Gauss' law and the conservation of global kinetic quantities. Another interesting observation is that already on lattices of rather small sizes, the scaling with lattice size appears.

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3.2. Infrared Instabilities in Nonabelian Gauge Theories

(C. Gong, S. G. Matinyan, B. Müller and A. Trayanov)

Recently the intriguing problem of the possibly large rate for electroweak baryon-number changing processes in high-energy collisions, originally based on the instanton approach [2], was considered from a completely different point of view, based on the classical analogue to high-energy particle collisions [3,4]. The hope was that one could observe the energy transfer from fast (high frequency) classical wave modes, presumably corresponding to initial states containing few high-energy particles, to slow (low frequency) classical wave modes, representing low-energy multi-particle final states. In contrast to negative results found in numerical studies in the (1+1)-dimensional abelian Higgs model [3] and the ϕ^4 -theory in (3+1) space-time dimensions [4], we found, by observing the instability of a standing wave in the SU(2) gauge theory, there exists such an energy transfer in non-abelian gauge theories.

The simplest ansatz satisfying the free Yang-Mills equation is an abelian standing wave

$$A_i^c(x, t) = \delta_{i3}\delta_{c3}A \cos k_0 x \cos \omega_0 t \quad (1)$$

with $k_0 = \omega_0$ assumed to be large. (We work in the temporal gauge $A_0^c = 0$.) To second order in A , we found perturbations $a_i^c(x, t)$ both transverse in color space and in real space contain instabilities. The strongest instability appears for a mode with $k_y = k_z = 0$, whose frequency contains an imaginary part $A/2\sqrt{2}$ in the limit of $k_x \gg A$. This is verified by a numerical study, the result of which is shown in Figure 1.

Initially the gauge configuration is an abelian standing wave with an amplitude corresponding to less than 5% of the maximal magnetic energy density on the 16^3 lattice. When a small *abelian* perturbation, restricted to the same direction in color space as the standing wave is applied, the field oscillations remain stable, as shown by the dashed line. When a general *nonabelian* perturbation pointing in a random color direction is added, the field oscillations develop a visible instability around time $t = 20$, as shown by the

solid curve. Quantitatively this instability can be characterized by the associated largest Lyapunov exponent λ_0 . In the continuum limit, we have

$$\lambda_0 \approx \frac{gA}{2}, \quad (2)$$

which is very close to the above perturbative result.

Fig. 1. The total electric energy as a function of time. The initial state is a standing wave with a small perturbation. The solid line shows the case when the perturbation is non-abelian, which results in the destruction of the plane wave beginning at $t \approx 20$. The dotted line is for an abelian perturbation, in which case the standing wave is stable. The computation is performed on a 16^3 lattice and energy per plaquette is about 0.16 in lattice units.

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3.3. Wilson Loop in Classical Lattice Gauge Theory (A. Trayanov)

For the pure gluon lattice gauge theory without quarks Wilson has proposed a gauge-invariant quantity that could serve as a non-local order parameter. The trace of a product of link variables around a closed loop is a gauge-invariant construction. Its expectation value is known as Wilson loop [1] and it can be expressed as

$$W(C) = \langle \text{tr} \prod_{ij \in C} U_{ij} \rangle \quad (1)$$

Here C is the contour of the loop and the group elements, U_{ij} are ordered as encountered in going around the path of the loop. The simplest non-trivial Wilson loop is the elementary plaquette on the lattice.

If a quark were to pass around the contour C , its wave function would pick up an internal symmetry rotation given by the product of the gauge link variables encountered. The Wilson loops essentially measures the response of the gauge fields to an external quark-like source passing around the perimeter.

It is expected that for large loops the loop expectation value will fall off exponentially with the area of the loop. The recent quantum study of the Wilson loops by Bali et al. [2] suggests that the string tension, which is the slope of this fall-off, is a function of the energy of the gauge fields. On the other hand perturbation theory arguments suggest that if a quantity is proportional to the dimensionless variable ag^2T , where a is the lattice constant, g is the gauge coupling and T is the temperature, then it will have the same value in the classical limits of the theory. The goal of this study is to find out whether this holds true for the string tension.

The method we use is similar to the one used for the time evolution study of the gauge fields [3]. The 3D space is discretized by a regular lattice, and the electric and magnetic gauge fields are defined on the links (ℓ), connecting two neighboring vertices (i, j), and on the elementary plaquettes p , forming the shortest closed path around a link, respectively. They are expressed in terms of SU(2) variables:

$$U_{ij} = \exp \left(-i \frac{1}{2} \tau^a A_\ell^a \right), \quad (2)$$

and

$$E_\ell^a = -i \frac{a}{g^2} \text{tr} [\tau^a \dot{U}_{ij} U_{ij}^{-1}], \quad \cos(\frac{1}{2} B_p) = \frac{1}{2} \text{tr} U_p. \quad (3)$$

Here U_p is the product of all four link variables on an elementary plaquette, and the dot denotes a time derivative.

The Hamiltonian of such system is the following:

$$H = \frac{g^2}{a} \left\{ \sum_{\ell} \frac{1}{2} E_\ell^a E_\ell^a + 4g^{-4} \sum_p [1 - \cos(\frac{1}{2} B_p)] \right\}. \quad (4)$$

The gauge fields evolve according to equations of motion. After choosing the initial conditions, i.e. the value of the total energy of the gauge fields, the chaotic dynamics of the fields leads fast to thermal equilibrium [3]. Then we periodically calculate the Wilson loop, W , (eq. [1]), for different loop sizes.

Currently we have developed a computer code which does the time evolution of the gauge fields and calculates the Wilson loops. The biggest lattice we can simulate on a single workstation is about 30^3 . Since we would like to study large Wilson loops we would like to explore even bigger lattices. Therefore we are implementing a parallel version of the code by splitting the lattice on different computing nodes by spatial decomposition. This will make the study of larger lattices a practical task.

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3.4. Wigner Functional Approach to Quantum Field Dynamics

(S. Mrówczyński and B. Müller)

Dynamical processes in relativistic quantum fields are usually described in terms of the particle excitations of these fields. However, there are situations where more appropriate degrees of freedom are rather the field amplitudes themselves and their conjugate momenta. This happens when as a first approximation the dynamics is well described by the classical field equations. A well known example here is the evolution of Higgs fields in the early universe, which acquire a finite vacuum expectation value when the temperature falls below the critical temperature for the symmetry breaking phase transition. The mechanism which drives the Higgs field from the symmetric phase, where the vacuum expectation value vanishes, to the asymmetric phase with a non-zero expectation value is essentially classical. Such phase transitions have attracted much attention in the context of inflationary cosmological models, models of galaxy formation, the creation of the cosmic baryon number asymmetry during the electroweak phase transition, and most recently due to the possibility of creating the so-called misaligned chiral condensates in ultrarelativistic heavy-ion collisions.

All these situations have in common that one must study the evolution of a quantum field far off thermal equilibrium, when the initial state of the quantum field is specified. Most theoretical approaches to this initial value problem for quantum fields have been based on dynamical equations for the field expectation values and the Gaussian fluctuations around those, assuming that fluctuations around the most probable path remain small. While this assumption may be correct in certain instances, it is a well known fact that fluctuations usually do not remain small in dynamical phase transitions, where domain formation and clustering are regularly occurring phenomena.

We have introduced a formalism which appears to be very convenient to study the evolution of quantum fields from an initial state and which goes beyond the classical description. The central object of our approach is the Wigner function, *Wigner functional* in fact, which is the quantum analog of the classical distribution function in a phase space. Here the phase space, however, is spanned not by the particle coordinates and momenta, but by the field amplitudes and their conjugate momenta. The great advantage of the Wigner function formalism is that, while being fully quantum mechanical, it remains close in its spirit to the classical description, and that the classical limit can always be easily obtained. On the other hand, the simultaneous presence of many "quasiclassical" field configurations can be described with ease, and the fluctuations around these do not have to remain Gaussian.

We have restricted our considerations to scalar fields with quartic self-interaction. After discussing the definition of the Wigner functional we have shown that the equations of motion have the familiar form of a transport equation in phase space with quantum corrections. We have then derived the Wigner functional for a free scalar field in thermal equilibrium. We have discussed how the classical phase space distribution is obtained in the high temperature limit, and we have obtained the two-point correlation function. We have analyzed the "rollover" of the scalar field in a second-order phase transition, where the power of the Wigner functional approach becomes apparent, allowing for the evolution of the quantum field along many simultaneous classical paths. We have briefly discussed the mean field and semiclassical approximation, and have suggested how to obtain numerical solutions to the initial value problem for quantum fields in more complex situations.

3.5. Wigner Function on Compact Spaces

(C. Gong and B. Müller)

One way to study the semiclassical limit of a field theory is via Wigner function approach. H-T. Elze and U. Heinz have worked with a Wigner operator, which for a scalar field assumes the form [1].

$$\hat{\mathcal{W}}(x, p) = \int \frac{d^3 y}{(2\pi\hbar)^3} e^{i p y / \hbar} \hat{\phi}^\dagger(x + y/2) \hat{\phi}(x - y/2), \quad (1)$$

where $\hat{\phi}(x)$ is the field operator. A second field Wigner function, which is a more direct generalization of the usual Wigner function, can be defined and is simplest within the lattice version of the field theory. Consider a scalar lattice field theory, with variables ϕ_s , where s denotes different lattice sites. The Wigner function is defined straightforwardly as

$$\mathcal{W}(\{\phi_s, \pi_s\}) = \int_{-\infty}^{+\infty} \left(\prod_s d\phi'_s \right) \langle \phi_s + \frac{1}{2}\phi'_s | \hat{\rho} | \phi_s - \frac{1}{2}\phi'_s \rangle e^{-i\pi_s \phi'_s}, \quad (2)$$

where $\hat{\rho}$ is the density operator and π_s is the conjugate momentum. But the above form cannot be directly implemented for a gauge theory. The difficulty stems from the fact that the variables in a lattice gauge theory live in a compact space; for example the link

variable U in the $SU(2)$ gauge theory is an element of $SU(2)$, the group manifold of which is a compact S^3 . We have solved this problem for the $U(1)$ theory, which in the one link limit is like a rotator. The solution for a non-abelian gauge theory is not known.

For a one link $U(1)$ gauge theory, the Wigner function is defined continuously on link variable U and discretely on eigenvalue m of the electric field E and the time evolution of this function is given by

$$i\partial_t \mathcal{W}_m(U) = m\hbar U \frac{\partial}{\partial U} \mathcal{W}_m(U) + \sum_n v_n(U) \mathcal{W}_{m-n}(U). \quad (3)$$

In the classical limit when m is large a continuous Wigner function can be defined and its time evolution is

$$i\partial_t \mathcal{W}(U, E) = \left(EU \frac{\partial}{\partial U} - \left(U \frac{\partial V(U)}{\partial U} \right) \frac{\partial}{\partial E} \right) \mathcal{W}(U, E). \quad (4)$$

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4. STUDIES IN QUANTUM FIELD THEORY

4.1. Renormalization Group and Universality

(S.-B. Liao [1])

The modification of fundamental laws of physics with changing observational length scale is the subject of renormalization group (RG). Through the RG flow equation, one may probe the dependence of the effective coupling constants on the characteristic length [2,3]. The otherwise complicated flow pattern becomes rather simple in the vicinity of the fixed points where the linearized RG flow along with scaling provide a recipe of classifying the coupling constants. The irrelevant coupling constants decrease as the scale is moved toward the infrared direction and therefore do not affect the physical contents in the infrared region. Within this regime where linearized RG becomes applicable and the usual concept of universality is recovered, physics is parameterized by the relevant and marginal coupling constants only.

In the realistic models we find several scaling regimes when the renormalized trajectory passes by different fixed points. There are fixed points for the Theory of Everything, GUT, Standard model, QCD and QED, to mention some of them. Although the true renormalized trajectory approaches all of them for certain values of the cut-off, it reaches the first one only. In fact, in the scaling regime of, say the fixed point of QCD, some of the interactions of the Standard Model generate non-renormalizable vertices in terms of the quark and gluon fields. These vertices deflect the renormalized trajectory from the fixed point as we move up in energy. In fact, when scaling is investigated at the other fixed points then the higher energy reactions always make these fixed points unstable in the ultraviolet direction [4].

We argue that the concept of universality is severely limited for models with multiple fixed points. As a demonstration, we derive an improved RG equation for the $O(N)$ symmetrical scalar field theory:

$$k\partial_k U_k(\Phi) = -\frac{k^4}{16\pi^2} \left\{ \ln \left[\frac{\tilde{Z}_{k,\ell} k^2 + U_k^{11}(\Phi)}{\tilde{Z}_{k,\ell} k^2 + U_k^{11}(0)} \right] + (N-1) \ln \left[\frac{\mathcal{Z}_{\parallel,\perp} \epsilon + \mathcal{U}_{\parallel}^{\epsilon\epsilon}(\oplus)}{\mathcal{Z}_{\parallel,\perp} \epsilon + \mathcal{U}_{\parallel}^{\epsilon\epsilon}(\rho)} \right] \right\}.$$

which accumulates the effects of the eliminated modes in a systematic way as the cut-off is lowered. Here the $\tilde{Z}$'s are the longitudinal and transverse wavefunction renormalization constants, and $U_k^{(ij)}$ denotes differentiation of U_k with respect to the i th and j th components of the scalar background field Φ . Our equation is superior compared with the usual one where only the contributions which are vanishing in the ultraviolet regime are retained.

The scaling properties around the infrared fixed point are expected to be uninteresting since less and less modes are kept in the system as the ultraviolet cut-off energy reaches zero. It is known that there is only one (completely trivial) relevant coupling constant

at this fixed point, the mass. But then we ignore the infrared divergences of the massless theories. The infrared divergences may in principle generate relevant operators as the ultraviolet cut-off approaches the infrared regime. Another class of models where the infrared scaling might be rather non-trivial is where a symmetry is broken spontaneously. Though the breaking of a symmetry influences the renormalized trajectory but the spontaneous breaking of the symmetry cannot be realized unless the system is infinite. Thus the symmetrical scaling of the ultraviolet becomes non-symmetrical only in the asymptotically low energy regime, in the expansion around the infinite system.

The form of the scaling operators of the infrared fixed points can be rather complicated in terms of the local field operators since the physics of the fixed point is essentially non-local. We recover classical physics at the fixed point for massless theories. There the soft particle emission allows the spread of the energy from the microscopical to macroscopical length scales. When the theory possesses a mass gap then the energy can not be distributed to arbitrary long distances and the infrared physics is still controlled by coherent quantum effects, e.g. superconductivity.

We found infrared singularities in the one-loop β function for the odd vertices of the $O(N)$ model when the first two orders in the gradient expansion of the renormalized lagrangian were retained. We know no more about the applicability of the gradient expansion than the trivial remark that it becomes more reliable at higher dimensions. Though one expects the gradient expansion to be reliable in the infrared regime, but it may not converge in determining the infrared scaling operators which might well be non-local. Thus we do not have an unambiguous support for the strong coupling infrared physics of the Goldstone modes. Nevertheless we think that the appearance of these infrared singularities is a serious issue. This is because our β function includes the complete one-loop result, in contrast to the usual computation where contributions which are vanishing at the ultraviolet fixed point are ignored. Thus our β function agrees with the ϵ expansion results in the vicinity of the ultraviolet fixed point. But in the infrared regime the contributions we retained play an important role and drive certain coupling constant to large values.

In summary, we find indications of the existence of relevant operators at the infrared fixed point, contrary to common expectations. This result makes the sufficiency of renormalizable coupling constants in parametrizing the long distance phenomena questionable.

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4.2. Renormalization Group and Gauge Invariant Renormalization (S.-B. Liao [1])

We derive a manifestly gauge invariant renormalized action for Yang-Mills theory using the Schwinger proper-time formalism. In our scheme the regularization is achieved by the smooth suppression of the large eigenvalues in the quantum fluctuation determinants.

The way how Λ and k enter gauge invariantly as the UV and IR regulator, respectively, can easily be seen by invoking the Schwinger proper-time formalism, in which the functional determinant of a general second-order operator $\mathcal{H}$ is represented by [2]:

$$\ln \left(\frac{\mathcal{H}}{\mathcal{H}_0} \right) = - \int_0^\infty \frac{dt}{t} (e^{-\mathcal{H}t} - e^{-\mathcal{H}_0 t}), \quad (1)$$

where the expression is normalized with respect to the corresponding operator $\mathcal{H}_0$ evaluated at vanishing field. If we introduce into the proper-time t integration a smearing function

$$\rho(z) = 1 - (1 + z)e^{-z}, \quad (2)$$

(2) then becomes, setting $z = \Lambda^2 t$,

$$- \int_0^\infty \frac{dt}{t} \rho(\Lambda^2 t) (e^{-\mathcal{H}t} - e^{-\mathcal{H}_0 t}) = \ln \left[\frac{\mathcal{H}}{\mathcal{H}_0} \times \frac{\mathcal{H}_0 + \Lambda^2}{\mathcal{H} + \Lambda^2} \right] - \frac{\Lambda^2 (\mathcal{H} - \mathcal{H}_0)}{(\Lambda^2 + \mathcal{H})(\Lambda^2 + \mathcal{H}_0)}, \quad (3)$$

which is the Pauli-Villars regularized version containing a UV cut-off.

Our approach makes it more feasible to implement Wilson's renormalization group method to Yang-Mills theory in a gauge symmetry preserving manner. In particular, we obtain a renormalization group flow equation which takes into consideration the contributions of higher dimensional operators and provides a systematic way of exploring the role of these operators as the strong coupling, infrared limit is approached [3].

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4.3. Renormalization Scheme Dependence of Radiative Symmetry Breaking (S.-B. Liao [1])

In the perturbative treatment of quantum field theory, one often encounters infrared (IR) divergences in evaluating Feynman graphs, especially when the system under investigation contains massless modes. Some of the renown examples in which the theories are plagued by IR singularities include QED and QCD due to the masslessness of the gauge particles. Such massless modes can also arise if the lagrangian of the theory possesses continuous global symmetry. Even though IR singularity reflects the long-range nature of

the interaction in the theory, it nevertheless plays an important role in justifying the perturbative approach to the high-energy phenomena. With the propagators of the massless theory behaving as p^{-2} , IR divergences occur in the vicinity of the mass-shell condition $p^2 = 0$. To generate a well-defined perturbation theory, one must therefore resort to some off-shell prescriptions.

Various schemes which have been suggested to bypass the complications of IR singularity in massless theories are listed below:

$$M \text{ scheme : } \begin{cases} \mu_R^2 = \frac{\partial^2 U_{M,k}}{\partial \Phi^2} \Big|_{\Phi=0,k=0} = 0 \\ \lambda_M = \frac{\partial^4 U_{M,k}}{\partial \Phi^4} \Big|_{\Phi=M,k=0} \end{cases}, \quad \tilde{M} \text{ scheme : } \begin{cases} \mu_{\tilde{M}}^2 = \frac{\partial^2 U_{\tilde{M}}}{\partial \Phi^2} \Big|_{\Phi=M,k=0} = 0 \\ \lambda_{\tilde{M}} = \frac{\partial^4 U_{\tilde{M}}}{\partial \Phi^4} \Big|_{\Phi=M,k=0} \end{cases}, \quad (1)$$

and

$$k \text{ scheme : } \begin{cases} \mu_R^2 = \frac{\partial^2 U_k}{\partial \Phi^2} \Big|_{\Phi=0,k=0} = 0 \\ \lambda_k = \frac{\partial^4 U_k}{\partial \Phi^4} \Big|_{\Phi=0,k} \end{cases}, \quad \tilde{k} \text{ scheme : } \begin{cases} \mu_{\tilde{k}}^2 = \frac{\partial^2 U_{\tilde{k}}}{\partial \Phi^2} \Big|_{\Phi=0,k} = 0 \\ \lambda_{\tilde{k}} = \frac{\partial^4 U_{\tilde{k}}}{\partial \Phi^4} \Big|_{\Phi=0,k} \end{cases},$$

where M and k are left as undetermined arbitrary non-vanishing parameters.

Upon careful comparisons, it is found that in spite of the fact that all conventional schemes yield the same β function for the coupling constant, there exists ambiguity in probing the symmetry of the vacuum in massless theory using the effective potentials derived from the above conditions. Different subtraction schemes in general lead to different forms of effective potentials, thereby making it difficult to decide on the symmetry of the ground state. Even though the M scheme [2] lends support to the occurrence of spontaneous symmetry breaking due to radiative corrections, the conclusion is refuted by the $\tilde{M}$ scheme. Such scheme dependence is exemplified in the well-known massless scalar QED and the two-dimensional Gross-Neveu models, in which spontaneous symmetry breaking is generally believed to take place. We argue that whether or not the vacuum of massless theory is broken by radiative correction still awaits for a more convincing method or experimental inputs to come.

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4.4. Effective Field Theory at Finite Temperature

(C. Gong, S.-B. Liao and B. Müller)

Blocking transformation is performed in quantum field theory at finite temperature for scalar field theory. Defining the coarse-grained blocked field as [1]:

$$\phi_k(\mathbf{x}) = T \int_0^\beta dy^0 \int d^3y \rho_k(\mathbf{x} - \mathbf{y}) \phi(\mathbf{y}), \quad \rho_k(\mathbf{x}) = \int_{|\mathbf{p}| < k} \frac{d^3\mathbf{p}}{(2\pi)^3} e^{i\mathbf{p} \cdot \mathbf{x}}, \quad (1)$$

the finite-temperature blocked potential is found to be [3]:

$$U_{\beta,k}(\Phi) = V(\Phi) + \frac{1}{2} \int_k^\infty \frac{d^3\mathbf{p}}{(2\pi)^3} \int_{-\infty}^\infty \frac{dp^0}{2\pi} \ln [p^2 + V''(\Phi)] \\ + T \int_k^\infty \frac{d^3\mathbf{p}}{(2\pi)^3} \ln [1 - e^{-\beta \sqrt{p^2 + V''(\Phi)}}]. \quad (2)$$

Since the manipulation above applies to arbitrary potential $V(\Phi)$, the renormalization group equation for the finite-temperature local potential $U_{\beta,k}(\Phi)$ can be obtained by differentiating (2) with respect to k :

$$k \frac{\partial U_{\beta,k}}{\partial k} = -\frac{k^3}{4\pi^2} \sqrt{k^2 + U''_{\beta,k}} - T \frac{k^3}{2\pi^2} \ln [1 - e^{-\beta \sqrt{k^2 + U''_{\beta,k}}}] . \quad (3)$$

The higher loop contributions are vanishing in the renormalization group equation (3) since they are suppressed by $\frac{\partial k}{k} \rightarrow 0$ [3]. To obtain a flow equation for the full theory, however, one needs to take into accounts the wave function renormalization constant as well as the higher order derivative terms in the blocked lagrangian. The simplifications, we believe, represent only small errors in the infrared limit of the four dimensional theories.

At high temperature the large Matsubara frequency suppresses the contribution of the non-static modes, hence terms with $n \neq 0$ have diminishing dependence on Φ . When the contribution of the non-static modes is altogether neglected then the static modes are described by a three dimensional theory with the potential $\tilde{U}_k = \beta U_{\beta,k}$. The natural field variable is $\tilde{\Phi} = \sqrt{\beta} \Phi$ in three dimensions and the renormalization group equation of the three-dimensional theory becomes:

$$k \frac{\partial \tilde{U}_k(\tilde{\Phi})}{\partial k} = -\frac{k^3}{4\pi^2} \ln \left[\frac{k^2 + \tilde{U}_k(\tilde{\Phi})}{k^2 + \tilde{U}_k(0)} \right], \quad (4)$$

where the dot denotes differentiation with respect to $\tilde{\Phi}$.

The power of finite-temperature RG flow equation 930 is that it already takes into accounts all the daisy and superdaisy diagrams. At every step when the momentum modes from each "thin shell" are being eliminated, the effective finite-temperature k -dependent effective vertex functions are generated, which are used subsequently as the

input parameters for the computation of $U_{\beta,k}(\Phi)$ at the next lower energy scale. Numerical works along this aspect are currently being pursued.

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4.5. Time-Dependent Variational Principle for ϕ^4 Field Theory (Chiyong Lin)

The study of time evolution of quantum systems is important in various areas of physics. In the high energy heavy-ion collisions, properties of hadronic matter can be explored through the globally off-equilibrium transient phenomena. On the other hand, the inflationary scenario of the early universe also involves the control of the time evolution of a driving scalar field. To explore these issues, the time-dependent variational method using a Gaussian trial wave functional has been utilized for ϕ^4 field theory. Although the formalism has been introduced for some time, the issue of renormalization of equations of motion was not completely discussed. We showed recently that the standard method for renormalization applies to these time-dependent equations. The crucial point is to use the hamiltonian nature of the variational principle. In particular we have considered small oscillations about equilibrium and shown that these give two meson modes of the theory. The two meson equation has a closed solution leading to a single bound state for attractive renormalised coupling and a complete form for the scattering amplitude in the continuum. This form is easily adapted to the usual running coupling constant in the two meson energy. We have also examined our equations for so called runaway solution where one of the physical parameters goes to infinity. Using the "potential" part of our variational hamiltonian we are able to show that conservation of energy prohibits any unphysical runaway.

4.6. Kinetic Approach to the Initial-Value Problem in Quantum Field Theory (Chiyong Lin)

Initial-value problems in quantum field theory such as non-equilibrium phenomena can be analyzed using, as the first crude attempt, the Gaussian approximation. However, as discussed by F.Cooper, S-Y Pi and P. Stancioff, this particular method suffers several shortcomings. Therefore, efforts have been devoted to finding an alternative scheme to improve the Gaussian approach. In our recent paper, a new method to solve such initial-value problems has been proposed. Using the time-dependent projection technique, we obtain a mean-field expansion, in which the lowest and higher-order terms correspond to, respectively, the Gaussian approximation and the dynamical correlation effects. This procedure was applied to ϕ^4 theory in $1+1$ dimension. We were able to compute the lowest two orders in the expansion for a spatially uniform system both analytically and numerically. Recently we have extended this technique to a non-uniform field configuration.

In addition, we have found some interesting aspects of the approach. It is possible to show that the differential equations obtained from this approximation represent a classical dynamical system. The system is hamiltonian for the mean-field approximation and "quasi-hamiltonian" when the kinetics of quantum correlations is incorporated to the dynamics. These properties are very important for example in the context of renormalizations for the equations of motion.

4.7. The Dirac-Coulomb Problem for the κ -Poincaré Quantum Group (B. Müller)

Quantum groups are a new symmetry structure recently introduced into physics. For semi-simple Lie algebras, the most important properties of the representation theory have been generalized to their quantum analogs and the theory and applications are by now well-developed. The situation is very different for non-semisimple groups where the quantum deformation structure is not unique. Some of the most fundamental, and interesting, symmetries in physics—most importantly the Poincaré and Galilei symmetries—fall into this category.

One approach to a deformed Poincaré algebra is to apply the contraction process to the standard q -deformation of the (anti) de Sitter algebra $so(3,2)$. Taking the limit of the de Sitter radius $R \rightarrow \infty$ with an accompanying limit of the deformation parameter q , such that $\lim(R \ln q) = \kappa^{-1}$, one obtains the κ -Poincaré quantum group. In contrast to the deformation parameter q , which is dimensionless, the parameter κ carries the units of a (fundamental?) length. The κ -Poincaré quantum group has two invariants, a quadratic invariant and a bi-quadratic invariant. The κ -Dirac equation has recently been found, and it factorizes the bi-quadratic invariant, quite unlike the usual case which factorizes the quadratic invariant.

It is a quite general question whether any of the quantum groups associated with the classical Lie groups are realized in nature. In particular, it would be interesting to probe whether any of the fundamental symmetries of nature, such as translational and Lorentz invariance, are represented at the quantum level precisely by Lie groups or rather by their quantum deformations. From the success of the standard description of these symmetries as Lie groups it is clear from the outset that the deformation parameter must be tiny, and that very precise tests are required to detect any deviation from the (undeformed) Lie group description. One of the most delicate tests of rotational symmetry in nature is provided by the near degeneracy of the $2s_{1/2}$ and $2p_{1/2}$ states in hydrogen, which is only broken by radiative corrections. This degeneracy has previously been used to test the dimensionality of space-time [1]. Here we study it in the context of the κ -Poincaré quantum group [2].

To obtain the κ -Dirac-Coulomb equation we have gauged the κ -Dirac equation and incorporated the Coulomb field. Expanding this equation in the (dimensional) parameter κ^{-1} , we apply it to the quantum relativistic hydrogen atom, with the aim of confronting any induced κ -Poincaré shift in the ground state energy level with the recent precision quantum optics measurements. Remarkably, we find that the first order effect *vanishes identically*,

while the second order perturbation contains singular terms rendering the equation ill-defined.

There is a rather speculative way out of this impasse. One of the motivations underlying the quantum group approach is that the quantum parameter may possibly serve as a convergence factor. The singular terms in the κ -Poincaré-Dirac-Coulomb problem enter through exponentials of the P_0 operator, which suggests that finite (time) displacements occur. Thus it is not unreasonable to hope that the κ -Poincaré length scale may somehow cut-off the singular terms in the Coulomb problem. For the singular second order perturbation term the cut-off enters logarithmically, and is thus relatively insensitive to the cut-off. Using the ground state second-order (cut-off) energy shift determines, in fact, a length-scale that is self-consistent (the cut-off scale in the singular integral is of the same order of magnitude as that scale determined by the cut-off energy shifts).

The ground-state second-order (cut-off) energy shift is found to have the order of magnitude:

$$\Delta E(1s_{1/2}) \approx (m\kappa)^2 \cdot (m\alpha^6).$$

Comparing this estimated shift, $\Delta E(1s_{1/2})$, with the observed experimental limits on the accuracy with which the $1s_{1/2}$ eigenenergy is known, we obtain:

$$(m\kappa)^2(m\alpha^6) \approx (m\alpha^2) \times 10^{-12}, \text{ or } m\kappa \approx 10^{-2}.$$

This implies that the κ -Poincaré length scale is: $\kappa \leq 10^{-13} \text{ cm}$ or smaller. We note that this self-consistent estimate of the κ -Poincaré length scale is both speculative and surprisingly inaccurate, despite the high experimental precision, in direct consequence of the remarkable property that the κ -Poincaré Dirac-Coulomb problem has no first-order correction to the discrete energy eigenstates.

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APPENDICES

IA. Published Articles:

1. Magnetic Screening in Thermal Gauge Theories, T. S. Biró and B. Müller, *Nucl. Phys.* **A561**, 477-500 (1993).
2. Strong and Electromagnetic Decays of the Baryon Decuplet, M. N. Butler, M. J. Savage and R. P. Springer, *Nucl. Phys.* **B399**, 69-88 (1993).
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6. Branching Processes and Multi-Particle Production, S. G. Matinyan and E. B. Prokhorenko, *Phys. Rev.* **D48**, 5127-5132 (1993).
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16. QCD Evolution Equations for High Energy Partons in Nuclear Matter, K. Geiger and B. Müller, *Phys. Rev. D* (in print).
17. Larger Domains of Disoriented Chiral Condensate by Annealing, S. Gavin and B. Müller, *Phys. Lett. B* (in print).
18. Thermal Photons and Dileptons from Non-Equilibrium Quark Gluon Plasma, M. T. Strickland, *Phys. Lett. B* (in print).

IB. Submitted Articles:

1. Pair Production in the Quantum Boltzmann Equation, J. Rau, submitted to *Phys. Rev. D*.
2. Universality and Renormalization Group, S.-B. Liao and J. Polonyi, submitted to *Phys. Rev. D*.
3. Multiple Strange Nuclear Systems, to appear in *Ann. Phys. BNL-49062* preprint.
4. Renormalization Group, Validity of One-loop Computation, and Triviality, S.-B. Liao and C. Gong, submitted to *Phys. Lett. B*.
5. Renormalization Group for Yang-Mills Theories, S.-B. Liao and J. Polonyi.
6. On Regularization Dependence in Infrared Theory, S.-B. Liao.
7. Pion Fusion in Peripheral pp Collisions, A. J. Schramm and B. Müller.

IC. Conference Reports:

1. D-Meson as a Probe of Early Parton Rescattering, X. N. Wang and B. Müller, *Nucl. Phys. A* **566**, 555c-558c (1994).
2. Strange Matter—A New Domain of Nuclear Physics, C. Greiner, D. Diener, J. Schaffner, and H. Stöcker, *Nucl. Phys. A* **566**, 157c-166c (1994).
3. Parton equilibration at RHIC and LHC, T. S. Biró, B. Müller, M. H. Thoma, and X. N. Wang, *Nucl. Phys. A* **566** (1994) 543c-546c.
4. Branching Processes and Multi-Particle Production, S. G. Matinyan, Proceedings of the Fifth Blois Workshop on Diffractive Processes at High Energy, June 1993, Providence, (World Scientific, 1994, p. 225).
5. Relativistic Heavy-Ion Collisions: Probing the Early Universe in the Laboratory, B. Müller, Proceedings of the Workshop on Pre-Equilibrium Parton Dynamics, LBL, Berkeley, August 1993, p. 1-39.

IIA. Invited Talks:

1. B. Müller, Chaos and Equilibration of Yang-Mills Fields on the Lattice, Workshop on Relativistic Heavy Ion Physics, Bergen, Norway, June 1993.
2. B. Müller, Physics of the Quark-Gluon Plasma, Gordon Research Conference on QCD in Nuclear Physics, Tilton, NH, June 26-30, 1993.
3. S. G. Matinyan, Negative Binomial Distribution in the High Energy Multi-particle Production, V Blois Workshop, Brown University, Providence, RI, June 1993.
4. C. Greiner, Strange Matter—A New Domain of Nuclear Physics, QM'93 Conference in Borlänge, Sweden, June 1993.
5. R. P. Springer, An Effective Theory for Nuclear Physics, NSF Summer School, Raleigh, NC, July 1993.
6. C. Greiner, Strange Matter—A New Domain of Nuclear Physics, Lecture series presented at the RHIC School '93, Brookhaven National Laboratory, August 1993.
7. B. Müller, Meson Production and Quark Gluon Plasma, Theory Workshop on Meson-Nucleus Dynamics at Intermediate and High Energies, Argonne National Laboratory, August 2-5, 1993.
8. B. Müller, Real Time Calculations in Lattice Gauge Theory, Conference on Strong Interactions at Finite Temperature, IPT Santa Barbara, August 16-21, 1993.
9. S. G. Matinyan, Instabilities in the Non-Abelian Gauge Theories and their Manifestation in High Energy Collisions, Institute of Theoretical Physics, University of Minnesota, October 1993.
10. C. Greiner, Strange Matter—A New Domain of Nuclear Physics, NATO Advanced Study Institute on Hot and Dense Nuclear Matter, Bodrum, Turkey, September 1993.
11. C. Greiner, Memory Effects in the Collision Term in Relativistic Heavy Ion Collisions, NATO Advanced Study Institute on Hot and Dense Nuclear Matter, Bodrum, Turkey, September 1993.
12. R. P. Springer, The Decuplet in Chiral Perturbation Theory, Lanczos Centenary Conference, North Carolina State University, Raleigh, NC, December 14, 1993.
13. C. Greiner, Strange Matter—A New Domain of Nuclear Physics, XVII Symposium on Nuclear Physics, Oaxtepec, Mexico, January 1994.
14. Sen-Ben Liao, Renormalization Group and Gauge Symmetry, Annual Strasbourg-Heidelberg Collaboration Meeting, February 29, 1994, Heidelberg, Germany.
15. B. Müller, Conditions for Quark-Gluon Plasma Formation, Rencontres de Moriond, Meribel, France, March 21-25, 1994.
16. B. Müller, QCD and Relativistic Heavy Ion Collisions, QCD Workshop, Madison, April 9-10, 1994.

IIB. Contributed Talks and Seminars:

1. J. Rau, The Apparent Irreversibility of Particle Creation, Institute for Nuclear Theory at the University of Washington, Seattle, May 1993.
2. B. Müller, Chaos in Nonabelian Gauge Theories, CEN Orsay, France, June 4, 1993.
3. C. Greiner, Memory Effects in the Collision Term in Relativistic HIC, GSI, Darmstadt, Germany, June 1993.
4. C. Gong, Classical Dynamics of Non-abelian Lattice Gauge Theory, Uehling Summer School at University of Washington, July 1, 1993.
5. C. Gong, Hamiltonian Dynamics and Chaotic Behavior of Yang-Mills Theories on a Lattice, Conference on Strong Interactions at Finite Temperatures program, University of California at Santa Barbara, October 21, 1993.
6. J. Rau, Ist die Erzeugung von Teilchen irreversibel?, Max-Planck-Institut fuer Kernphysik, Heidelberg, Germany, November 4, 1993.
7. J. Rau, The Projection Method: a Universal Link between Microdynamics and Macrodynamics, Universite Louis Pasteur/Centre de Recherches Nucleaires, Strasbourg, France, January 26, 1994.
8. J. Rau, Pair Production in the Quantum Boltzmann Equation, Universite Louis Pasteur/Centre de Recherches Nucleaires, Strasbourg, France, January 27, 1994.
9. B. Müller, Relativistic Heavy Ion Collisions: Probing the Early Universe in the Laboratory, Colloquium, University of Minnesota, February 2, 1994.
10. S. G. Matinyan, Dynamical Chaos of Non-Abelian Gauge Fields, Brown University, March 1994.
11. R. P. Springer, Chiral Lagrangians with Baryons, Colloquium, Clemson University, April 7, 1994.
12. C.-Y. Lin, Time-Dependent Variational Principle for ϕ^4 Field Theory: RPA Approximation and Renormalization, Yale University, April 25, 1994.
13. J. Rau, Pair Production in the Quantum Boltzmann Equation, Columbia University, April 26, 1994.
14. R. P. Springer, Issues in Chiral Perturbation Theory, University of Maryland, April 29, 1994.
15. Sen-Ben Liao, Renormalization Group Approach to Quantum Field Theory, Academia Sinica, Taipei, May 7, 1994.
16. Sen-Ben Liao, Renormalization Group and Universality, National Taiwan University, May 9, 1994.
17. J. Rau, Pair Production in the Quantum Boltzmann Equation, University of Colorado at Boulder, May 12, 1994.

18. Sen-Ben Liao, Renormalization Group and Gauge Symmetry, National Taiwan University, May 12, 1994.

III. Triangle Nuclear Theory Colloquium Series

- 9/7 Ian Kogan (Princeton)
Disoriented chiral condensates
- 9/14 Wick Haxton (Seattle)
Electric dipole moment constraints on T-even, P-odd nucleon-nucleon interactions
- 9/21 Tom Cohen (University of Maryland)
Just how strange is the nucleon?
- 9/28 Dieter Hartmann (Clemson University)
Neutron stars and gamma-ray bursts
- 10/5 Pierre Sikivie (University of Florida)
Dark matter axions
- 10/12 Benjamin Grinstein (SSC)
Form Factors for B-meson Decays
- 10/26 Mike Musolf (ODU/CEBAF)
Searching for Strangeness at CEBAF
- 11/2 John Clark (Washington University)
Nucleonic superfluidity
- 11/9 Frank Tabakin (University of Pittsburgh)
Spin observables in vector meson photoproduction
- 11/16 Andrew Sandorfi (Brookhaven)
How round is the proton? - Physics with LEGS
- 11/23 Steve Wallace (University of Maryland)
Electromagnetic scattering from a relativistic bound state: the deuteron
- 11/30 Robert Malaney (University of Toronto)
Neutrino lasers and mixed dark matter
- 12/14 Mark Strikman (Penn State University)
Color fluctuations and collisions of hadrons with nuclei
- 1/18 Bruce Barrett (University of Arizona)
The interacting boson model for nuclei, atoms and molecules
- 1/25 Ubirajara van Kolck (University of Washington)
Chiral Lagrangians and nuclear forces
- 2/1 Stuart Pittel (Bartol Res. Institute)
Baryon mappings for deriving nuclei from constituent quarks

- 2/8 Henry Weller (TUNL/Duke)
p + Li and p+d capture at energies below 100 keV
- 2/22 Don Lehman (George Washington University)
Continuum Faddeev calculations for electromagnetic processes
- 3/1 Jim Napolitano (Rensselaer Polytechnic Institute)
Explorations of the baryon spectrum: New twists on an old subject
- 3/8 Vladimir Braun (MPI München)
Nonlocal operator product expansion and higher twist effects in quantum chromodynamics
- 3/15 Jitendra Parikh (Ahmedabad)
The chaotic quark-gluon plasma
- 3/22 Robert Perry (Ohio State University)
Light-front QCD and the renormalization group
- 3/29 Alan White (Argonne National Lab)
The QCD pomeron and small-x physics
- 4/5 David Kosower (Saclay)
String theory applied to QCD
- 4/12 Gerry Brown (Stony Brook)
Where the nuclear pions are
- 4/26 Ian Towner (Chalk River)
Superallowed Fermi beta decay
- 5/17 Eric Swanson (MIT)
Distinguishing stray decay models for mesons

IV. Theory Seminars at Duke

- 4/19 Heribert Weigert (Universität Regensburg)
Radial gauges for kinetic theories of "hot" QCD
- 5/7 Sergei Matinyan (Yerevan/Duke University)
Masses of top quark and Higgs bosons and CP violation
- 6/1 Pierre van Baal (Lorentz-Instituut, Univ. Leiden)
Glueball spectroscopy in finite volumes
- 9/8 Ian Kogan (Princeton University)
Axions, monopoles and vortices: from universe to hetero-structures
- 9/15 John Board (Electrical Engineering, Duke Univ.)
Recent approaches to the N-body problem of electrostatics and gravitation

- 9/20 Chengqian Gong (Duke University)
Hamiltonian dynamics of Yang-Mills fields on a lattice
- 9/22 Carsten Greiner (Duke University)
What's new in hypernuclei
- 9/29 Richard G. Palmer (Duke University)
Statistical mechanics of genetic algorithms
- 10/6 Roxanne Springer (Duke University)
The uses and abuses of chiral perturbation theory
- 10/13 John Risley (N.C. State University)
The role of computers in physics education
- 10/27 Sen-Ben Liao (Duke University)
Can a particle generate its own mass? Blocking transformations
and domain formation in quantum field theory
- 10/27 Neil Turok (Princeton University)
Large scale structure formation in the universe
- 10/28 Takayuki Kubo (Tokyo University)
Beyond chaos in fermion systems
- 11/1 John W. Clark (Washington University)
Neural networks that learn to predict probabilities:
An application to nuclear science
- 11/2 Steffen A. Bass (University of Frankfurt, Germany)
Collective effects of mesons in nuclear collisions
- 11/8 Fred Olness (Southern Methodist University)
Not so heavy flavor production: can QCD predict it?
- 11/10 Chengqian Gong (Duke University)
Infrared instabilities of nonabelian waves
- 11/15 A. G. Sedrakian (Yerevan Physics Institute)
Three dimensional Ising Model as a string theory
- 11/17 Steve Reynolds (North Carolina State University)
Computational astrophysics
- 11/18 Sergei Matinyan (Yerevan/Duke University)
Regge poles for beginners (Series of Seven Lectures)
- 11/24 Chi-Yong Lin (Duke University)
Time-dependent variational principle for quantum field theory
- 12/15 Mark Strikmann (Penn. State University)
Color transparency

- 1/10 Berndt Müller (Duke University)
Disordered chiral condensates: one year later
- 1/12 Michael Strickland (Duke University)
Photons and leptons from the quark gluon plasma
- 1/17 Roxanne Springer (Duke University)
CP violation
- 1/26 Ubirajara van Kolck (University of Washington)
Chiral symmetry and isospin symmetry breaking
- 2/9 David Morrison (Duke University)
Space-time topology change in string theory (Part 1)
- 2/16 David Morrison (Duke University)
Space-time topology change in string theory (Part 2)
- 2/18 Viktor Flambaum (University of New South Wales, Australia)
Statistical enhancement of weak interactions and quantum chaos in atoms and nuclei
- 2/23 Senben Liao (Duke University)
Renormalization group and gauge invariant regularization
- 3/2 Andreas Schäfer (University of Frankfurt, Germany)
Some current issues in small- x Physics
- 4/6 Stanisław Mrówczyński (University of Warszawa)
Statistical aspects of particle production in relativistic heavy ion collisions
- 4/13 Eugen Merzbacher (University of North Carolina)
The inexorable growth of the physical review: A multidimensional problem
- 5/3 Jochen Rau (MPI Heidelberg/Duke University)
The "black box" approach to universal conductance fluctuations
- 5/4 Senben Liao (Duke University)
Renormalization group and universality
- 5/15 Vyacheslav Belyaev (CEBAF)
 $\pi - A_1$ electromagnetic form factors and light-cone QCD sum rules