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DYNAMICS OF THE CHIRAL TRANSITION

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ABSTRACT

Measurements of disoriented chiral condensates (DCC) in heavy ion collisions at RHIC can yield fundamental information on the nature of the QCD phase transition. I review theoretical efforts to understand DCC formation and present work in progress on possible experimental ramifications.

Invited talk presented at the 5th Conference on
The Intersections of Particle and Nuclear Physics
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Abstract

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The order of the chiral restoration phase transition in high temperature QCD is currently unknown. If the transition is second order, Rajagopal and Wilczek (RW) propose that the nonequilibrium dynamics in ion-ion collisions at RHIC can generate transient domains in which a macroscopic pion field develops [1]. These domains are called disoriented chiral condensates (DCC). As Taylor discussed at this meeting, DCC's can lead to fluctuations in the charged and neutral pion spectra. In ion-ion collisions, the ability of experimenters to identify DCC's amidst the background produced by particle cascading critically depends on the domains' size and energy content. In this talk, I explore the RW mechanism for domain formation to determine the size of domains in nuclear collisions, reporting on work with Gocksch and Pisarski [2] and Müller [3]. I then briefly speculate on phenomenological consequences of DCC formation.

Equilibrium high temperature QCD manifests a chiral symmetry if the light up and down quarks are taken to be massless. However, a phase transition occurs at the critical temperature $T_c \sim 140$ MeV in which chiral symmetry is broken by the formation of a scalar $\langle \bar{q}q \rangle$ condensate. In the real world chiral symmetry is explicitly broken at all temperatures by the few-MeV current quark masses. In that case the nature of the transition has yet to be established [4].

Rajagopal and Wilczek pointed out that the chiral condensate can become temporarily disoriented in the nonequilibrium conditions encountered in heavy ion collisions. Near T_c , the approximate chiral symmetry implies that the scalar condensate is nearly equivalent to a pion-like pseudoscalar isovector condensate $\sim \langle \bar{q} \gamma_5 \vec{\tau} q \rangle$, where $\vec{\tau}$ are Pauli isospin matrices. Consequently, domains containing a macroscopic pion field can appear as the temperature drops below T_c . Such domains will eventually disappear as the system evolves towards the true vacuum in which only the scalar condensate is nonzero. In the finite-sized heavy ion system, however, the evolving domains can radiate pions preferentially according to their isospin content.

The excitation of a pion field over a substantial fraction of the collision volume can result in novel fluctuations in the number of neutral and charged pions. [See Taylor's talk in these proceedings.] Consider a single ideal domain in which the pion field is oscillating along a fixed isospin direction $\vec{\pi}$. Pion production from this domain is proportional to

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Equilibrium high temperature QCD manifests a chiral symmetry if the light up and down quarks are taken to be massless. However, a phase transition occurs at the critical temperature $T_c \sim 140$ MeV in which chiral symmetry is broken by the formation of a scalar $\langle \bar{q}q \rangle$ condensate. In the real world chiral symmetry is explicitly broken at all temperatures by the few-MeV current quark masses. In that case the nature of the transition has yet to be established [4].

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The excitation of a pion field over a substantial fraction of the collision volume can result in novel fluctuations in the number of neutral and charged pions. [See Taylor's talk in these proceedings.] Consider a single ideal domain in which the pion field is oscillating along a fixed isospin direction $\vec{\pi}$. Pion production from this domain is proportional to

the square of the field strength. All orientations of $\vec{\pi}$ are equally likely, so that the probability that a domain radiates the fraction of neutral pions f is

$$P(f) \propto d(\cos \theta)/df \propto f^{-1/2} \quad \text{where} \quad f = |\pi^0|^2 / \sum_{0,+,-} |\pi^a|^2 \equiv \cos^2 \theta. \quad (1)$$

It is therefore most probable for a single domain to emit more charged pions than neutral.

The detectability of domains in a nuclear collision depends on their size [2]. If the interaction volume in a nuclear collision is dominated by a single domain, then DCC formation can lead to a measurable isospin asymmetry. If a nuclear collision instead produces many uncorrelated small domains, the spectrum of fluctuations would be gaussian, as expected, *e.g.*, if the pions were produced independently by cascading.

To develop a more concrete picture of how the chiral condensate evolves, we use the linear sigma model, in which the pion field is coupled to a scalar σ field that characterizes the scalar condensate [1]. The fields interact through a potential $V = \lambda(\vec{\pi}^2 + \sigma^2 - v^2)^2/4 - H\sigma$. Chiral symmetry transformations correspond to $O(4)$ rotations of the vector field $\Phi = (\sigma, \vec{\pi})$. For $H = 0$ this symmetry is spontaneously broken at zero temperature, with $\langle \sigma \rangle = f_\pi = 93$ MeV. The external field H ($\propto$ the light quark masses) breaks the $O(4)$ symmetry explicitly and gives the pions a small mass $m_\pi = \sqrt{H/f_\pi} = 140$ MeV. In comparison, the sigma mass $m_\sigma \sim \sqrt{2\lambda}v$ is quite large, perhaps ~ 600 MeV.

I take this model with the parameters fixed at $T = 0$ as the basis for a Ginzburg-Landau description of the chiral transition [1]. For $H = 0$, the transition would then be second order with an $O(4)$ order parameter Φ . DCC domains would then be well-defined entities in the thermodynamic limit. While there is strictly no transition for $H \neq 0$, *transient* domains can form as follows.

In the idealized ‘quench’ scenario proposed by Rajagopal and Wilczek, one assumes that the system initially has $\Phi = 0$, as appropriate at high temperature, and follows the development of the system using $T = 0$ equations of motion derived from V . The initial state of the system is clearly unstable. The system “rolls down” from the unstable local maximum of $V(\Phi)$ towards the nearly stable values with $|\Phi| = v$ (the symmetry breaking term $-H\sigma$ is relatively small). 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 continue until interactions eventually damp the motion.

To be more concrete, the linearized equations of motion for the Fourier components of the pion field are:

$$\frac{d^2}{dt^2} \vec{\pi}_{\vec{k}} = \left\{ \tau_R^{-2} - k^2 \right\} \vec{\pi}_{\vec{k}} \quad \text{where} \quad \tau_R = \{\lambda v^2\}^{-1/2} \sim \sqrt{2}/m_\sigma. \quad (2)$$

Field configurations with $\langle \Phi \rangle = 0$ and momentum $k < \sqrt{\lambda}v$ are unstable and grow exponentially; modes with higher momenta do not grow. The $k = 0$ mode grows the fastest, with a time scale τ_R . The unstable modes grow for a time of order τ_R until Φ^2 approaches v^2 , after which the system oscillates about $\sigma = f_\pi$.

Significantly, the time during which growth can occur is quite short, $\tau_R \sim 0.5$ fm for $m_\sigma \sim 600$ MeV. Rajagopal and Wilczek found that the power $\propto \langle \pi^a_{\vec{k}} \pi^a_{-\vec{k}} \rangle$ in the

low momentum pion modes grows when the exact classical equations of motion are integrated, demonstrating that domains indeed form. [Here $\langle \dots \rangle$ represents an average over the initial of the fields, which are taken to have a thermal spectrum.] What is not clear, however, is whether the time during which the system is unstable is sufficient for truly large domains to form.

Gocksch, Pisarski and I [2] studied the domain size by numerically integrating the equations of motion in the quench scenario and extracting the spatial correlation function:

$$\langle \pi(\vec{x}, t) \pi(0, t) \rangle \propto \int d^3k \langle \pi_{-\vec{k}} \pi_{\vec{k}} \rangle e^{-i\vec{k} \cdot \vec{x}}, \quad (3)$$

where isospin labels are implicit. This correlator contains size information because $\langle \pi(\vec{x}, t) \pi(0, t) \rangle \approx \langle \pi(0, t) \rangle^2$ holds inside the domain. The numerical results in Ref. [2] can be understood by taking $\pi_k(t)$ from the linearized equation (2) and assuming that the fluctuations are initially thermal. A saddle-point integration yields $\langle \pi(\vec{x}) \pi(0) \rangle \propto e^{-x^2/8\tau_R t}$. The domain size, $R_D \sim 2\tau_R \propto m_\sigma^{-1}$, is set by the sigma mass. For $m_\sigma = 600$ MeV, we find that domains are essentially pion sized. Numerical simulations confirm this result.

The lesson drawn from the quench scenario is that large nucleus-sized domains can occur only if the sigma mass is small [2]. However, it is quite possible that the effective m_σ is reduced in the high energy density heavy ion environment. In fact, at a truly second order phase transition we expect m_σ — the inverse correlation length — to vanish! In that case, the parameters of the effective potential for our Ginzburg-Landau model cannot be *naively* taken from zero-temperature physics as we have done.

To explore the role of the medium in domain formation, Müller and I studied the dynamical evolution of the condensate in the presence of a nonequilibrium bath of quasiparticles [3]. We find that larger domains are possible due to two effects, *i*) the reduction of m_σ and *ii*) the “annealing” of the system, rather than quenching, by the slowly expanding heavy ion system. [Our use of the word “annealing” is meant to distinguish our scenario from RW’s “quench” — the similarity of these terms to those used in condensed matter physics is in some respects misleading.]

I illustrate the role of mass reduction by considering a system near equilibrium. The system is described by an effective potential V_{eff} that has the behavior shown in Fig. 1. In the Hartree mean-field approximation, the change of $V_{\text{eff}}(\Phi)$ as a function of temperature is determined essentially by the coefficient of Φ^2 , *i.e.* the mass term. Taking $H = 0$, one obtains a linearized equation analogous to (2) that implies a growth rate $\tau_{R,\text{eff}}^{-1} \propto m_\sigma^{\text{eff}} \propto (T_c^2 - T^2)^{1/2}$. Growth is very slow near the $T_c \approx \sqrt{2}v$, prolonging the time over which the system is unstable. Only as the potential approaches its free space shape does the roll-down become rapid.

The annealing scenario describes the nonequilibrium evolution of the system in a regime where the time scale τ_{ev} for the evolution of V_{eff} is much longer than τ_R . One expects such a slow cooling scenario to further enhance the domain size — in metallurgy, the more rapidly one quenches, the smaller are the resulting crystals. The evolution of V_{eff} is determined by the expansion of the surrounding quasiparticles, which constitute a time-varying nonequilibrium heat bath. The annealing regime is relevant at RHIC energies because the temperature falls to $T_c \approx 140$ MeV only at very late times, perhaps $\tau_c \sim 10 - 20$ fm. The expansion time τ_{ev} is then $\sim \tau_c \gg \tau_R \sim 0.5$ fm.

Figure 1: The effective potential for various temperatures.

Müller and I explored the evolution in the annealing regime using a nonequilibrium Hartree-like approximation. [See Ref. [3] for details.] We indeed find that annealing prolongs the time during which the system is unstable compared to a quench. Domain sizes up to 7 fm are possible if we assume that $m_\sigma^{\text{eff}}(T_c) = 0$ (i.e., a second order transition). More conservatively, we find sizes of 3–4 fm for $m_\sigma^{\text{eff}}(T_c) = 300$ MeV. Such large masses are possible, e.g., in the large N limit of the $O(N)$ model studied by Kluger in these proceedings. However, QCD may be far richer than the large N limit implies.

Suppose that a nuclear collision produces a large domain, either in an typical event due to the annealing dynamics, or due to rare fluctuation. How can we tell? The number of pions is not large, even ideally. The amount of energy available for pion production per unit volume is the potential difference $\Delta V \sim 60 \text{ MeV fm}^{-3}$ between $\Phi = 0$ and $\Phi = (f_\pi, \vec{0})$. If all of the energy in a domain of size $R_D \sim 3 - 6$ fm goes into pion production, one expects $N_{\text{dcc}} \sim \Delta V R_D^3 / m_\pi \sim 30 - 300$ pions. In comparison, ordinary cascading in a central Au+Au collision at RHIC can produce 1000 pions *per unit rapidity*; this constitutes a background to the DCC signal.

The source of pions from this domain is coherent and, therefore, concentrated at low transverse momenta. Horn and Silver, Gyulassy *et al.*, and Amado *et al.* have developed a coherent state formalism applicable to DCC production [5]. The amplitude for emitting a single neutral pion from a space time point (x_j, t_j) within a domain is real and $\propto \pi^0(x_j) \propto \exp\{-x_j^2/2R_D^2\}$. The production rate from a single domain is

$$E(dN/d^3p)_{\text{dcc}} \propto \left| \sum_j \pi^0(x_j, t_j) e^{-ip \cdot x_j + iE_p t_j} \right|^2 \sim e^{-p^2 R_D^2}, \quad (4)$$

neglecting an energy-dependent factor that is slowly varying for small momenta. The rate is largest for $p_T < R_D^{-1}$.

Experimentally, one must search for DCC's on an event-by-event basis, looking for significant fluctuations in the neutral relative to the charged pions. The DCC signal (4) appears atop an isospin symmetric background. Assuming that all the DCC pions are neutral and taking the background produced by the lund/fritiof event generator for central Au+Au, I obtain the spectra in Fig. 2. A signal that is significant compared

Figure 2: Neutral pion production from a single large DCC in Au+Au at RHIC.

to statistical fluctuations in the background emerges at low p_T for a domain of size $R_D > 3$.

How do we search for such domains, and what do they tell us? A domain producing a neutral fraction f introduces a contribution $fE(dN/d^3p)_{\text{dcc}}$ to the π^0 spectrum, with a corresponding contribution $\{(1-f)/2\}E(dN/d^3p)_{\text{dcc}}$ to the π^+ and π^- spectra. The distribution of f in a sample of large-domain events is given by (1). These properties hold for any coherent source [6]. If the coherence is produced by a nearly-second-order phase transition, then large-domain events would constitute a large fraction of central collisions and the pion excess would be independent of the background event structure.

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