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EXPLORING THE CHIRAL PHASE TRANSITION IN HIGH-ENERGY COLLISIONS

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Chiral symmetry plays an important role in nuclear physics but its manifestations have so far been testable only in the normal phase where the symmetry is spontaneously broken. More global explorations can be made by means of high-energy nuclear collisions, since a partial restoration of chiral symmetry is expected to occur within the highly agitated collision zone. Our current understanding of the chiral phase diagram will be reviewed, the special character of the induced non-equilibrium relaxation dynamics will be elucidated, and two areas of intensified future research will be highlighted.

1 Introduction

As the new century dawns, the field of nuclear physics encompasses a broad spectrum of research frontiers at which many promising research issues present themselves. Chiral symmetry constitutes a central concept in modern nuclear physics[1] and the present discussion focusses on the prospects for exploring its manifestations over a wide range of environments by means of high-energy nuclear collisions.

The strong interaction possesses chiral symmetry only to the degree that the basic quark masses can be neglected. Moreover, the self-interaction of the gluon fields encourages the spontaneous breaking of the symmetry in systems at low excitations, causing the familar vacuum to be a chiral condensate with a finite expectation value. Although chiral symmetry has proven to provide a successful framework for understanding many features in hadronic physics, the confrontation with experimental data has been primarily limited to the strongly broken phase.

In the highly excited environments produced in high-energy collisions, the chiral condensate is significantly weakened and the system may (briefly) achieve approximate chiral symmetry. The subsequent non-equilibrium relaxation process may then produce observable signals that could be exploited to test our global understanding of chiral symmetry in strong interactions[2,3,4].

This presentation reviews our current expectations regarding this novel phenomenon and identifies two important topics for intensified research, namely the development of a practical treatment of real-time non-equilibrium quantum-field dynamics in a non-uniform environment and the extension of the models from the *ud* SU(2) sector to the *uds* SU(3) sector of the strong interactions.

2 Equilibrium Properties

A simple theoretical tool for the global exploration of chiral symmetry in baryon-free systems is provided by the linear σ model. It describes the O(4) chiral field $\boldsymbol{\phi}(\mathbf{r},t)=(\sigma,\boldsymbol{\pi})$ by means of a simple effective quartic interaction,

$$\mathcal{L} \;=\; \tfrac{1}{2}\partial_\mu\boldsymbol{\phi}\circ\,\partial^\mu\boldsymbol{\phi} \;-\; \tfrac{\lambda}{4}(\boldsymbol{\phi}\circ\boldsymbol{\phi}-v^2)^2 \;+\; H\sigma \quad , \;\boldsymbol{\phi}\circ\boldsymbol{\phi} \;=\; \sigma(\mathbf{r})^2+\boldsymbol{\pi}(\mathbf{r})\cdot\boldsymbol{\pi}(\mathbf{r})\;. \quad (1)$$

The parameters, λ, v, H, may be fixed to reproduce the pion decay constant f_π, the free pion mass m_π, and the mass of the schematic σ meson, m_σ.

Useful insight can be gained by considering macroscopically uniform matter. The chiral field can then be decomposed, $\boldsymbol{\phi}(\mathbf{r},t)=\underline{\boldsymbol{\phi}}(t)+\boldsymbol{\delta\phi}(\mathbf{r},t)$, where $\boldsymbol{\delta\phi}$ represents quasi-particle agitations relative to the (false) vacuum characterized by the O(4) chiral order parameter $\underline{\boldsymbol{\phi}}$. Figure 1 (*left*) shows the associated free energy density $F_T=V_T(\phi_0,\chi_0)-TS_T(\phi_0)$ along the σ axis, with $\sigma_0=\phi_0\cos\chi_0$. At each temperature T, the minimum in F_T represents the equilibrium value of the order parameter and it shrinks steadily as T is increased, thus producing a *crossover* from the strongly broken to the approximately symmetric chiral phase.[?]

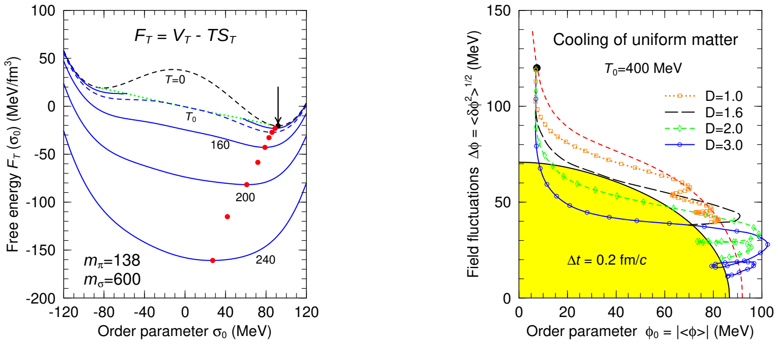


Figure 1: *Left:* The free energy density F_T calculated with the semi-classical mean-field approximation to the linear σ model, as a function of the order parameter σ_0 for various values T of the quasi-particle gas. At each temperature, the solid dots indicate the equilibria. For $T<T_0=\sqrt{2}v$, the order parameter must exceed a certain minimum value before all quasiparticle modes are stable; the corresponding end points are connected by the dotted curve. The top dashed curve is the bare potential obtained for T=0. *Right:* The combined dynamical evolution of the order parameter $\sigma_0=\langle\sigma\rangle$ and the field fluctuations $\langle\boldsymbol{\delta\phi}^2\rangle^{1/2}$. The dashed curve connects the equilibria from T=0 to above 500 MeV and the unstable region within which $\mu_\pi^2<0$ is shown by the shaded region. Each system has been prepared in thermal equilibrium at $T_0=400$ MeV and then subjected to a Rayleigh cooling, emulating a uniform expansion in D dimensions.

3 DCC Dynamics

If the field fluctuation is employed as a measure of the agitation, as in fig. 1 (*right*), the resulting diagram applies to any state of the system and is thus suitable for displaying non-equilibrium dynamics. In particular, it is possible to display the trajectories that result from the application of a Rayleigh cooling to initially hot chiral matter in equilibrium,

$$[\Box + \lambda(\boldsymbol{\phi}\circ\boldsymbol{\phi} - v^2)]\boldsymbol{\phi}(\mathbf{r},t) \;=\; H\boldsymbol{e}_\sigma - \frac{D}{t}\frac{\partial}{\partial t}\boldsymbol{\phi}\ , \tag{2}$$

which emulates a Bjorken-type scaling expansion in D dimensions.[6]

Two general features are apparent: 1) only very rapid cooling rates lead into the supercritical region where spontaneous pion creation occurs, thus suggesting that this interesting *quench* scenario may be dynamically unreachable; 2) for any cooling rate, the order parameter ultimately settles into rather regular damped oscillations around the true vacuum which may be expected to produce *parametric amplification* of pionic modes with frequencies $\omega_\pi \approx \frac{1}{2}m_\sigma$.

More refined scenarios considering rods endowed with a longitudinal Bjorken scaling expansion confirm the above expectations, as is illustrated in fig. 2 for either the dynamical trajectory (*left*) or the effective pion mass (*right*).[7]

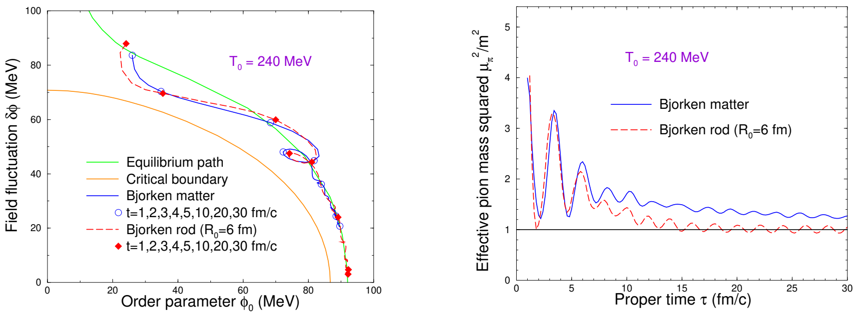


Figure 2: *Left:* Phase evolution of the interior of a Bjorken rod prepared with a bulk temperature of T_0=240 MeV and a radius of R_0=6 fm and the corresponding evolution of Bjorken matter (obtained in the limit of a very thick Bjorken rod, $R_0 \to \infty$) In order to be able to display an arbitrary non-equilibrium scenario, the average field fluctuation $\langle\delta\phi^2\rangle^{1/2} = \langle\delta\sigma^2 + \delta\pi^2\rangle^{1/2}$ has been used in place of the temperature on the vertical axis. Also shown are the equilibrium path and the boundary of the supercritical region. The rod has been probed over a hollow cylindrical volume, $1<\rho<3$ (fm). The paths are marked at successive proper times τ=1,2,3,4,5,10,20,30 fm/c. The two evolutions are similar early on, but the rod then relaxes significantly faster due to the generated transverse expansion. *Right:* The corresponding evolutions of the squared effective pion mass μ_π^2.

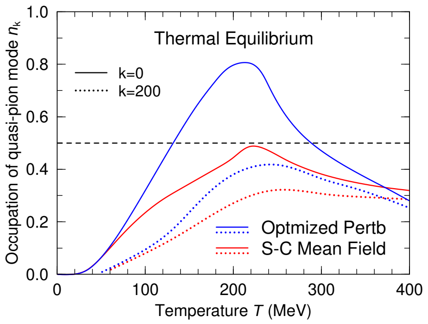
4 Vacuum Fluctuations

A full quantum-field treatment of the chiral dynamics is beyond current reach and the dynamical simulation studies have employed classical fields. Although much valuable insight can been gained in this manner, it is important to recognize that such treatments are not quantitatively reliable. This is perhaps best illustrated by considering a free pionic mode with a time-dependent mass, $[\Box + \omega_k^2(t)]\phi_{\mathbf{k}}(t) = 0$, where $\omega_k^2(t) = k^2 + \mu^2(t)$. If the initial occupancy of the mode is $n_{\mathbf{k}}^{\rm init}$, then the (phase-averaged) final occupancy becomes[8]

$$n_{\mathbf{k}}^{\rm final} \;=\; X_k \left[n_{\mathbf{k}}^{\rm init} \;+\; \tfrac{1}{2}\right] \;-\; \tfrac{1}{2} \;>\; X_k n_{\mathbf{k}}^{\rm init} \;, \tag{3}$$

which exceeds the classical result $X_k n_{\mathbf{k}}^{\rm init}$, since the amplification coefficient X_k is generally larger than unity. The expression brings out the fact that the quantum fluctuations and the statistical fluctuations combine in a democratic fashion as seeds for the parametric amplification. A quantitative impression of their relative importance can be gained from fig. 3 (*left*). It is evident that the thermal occupancies are never large compared to one half. Consequently, the vacuum fluctuations are never negligible.

Efforts to develop a suitable quantum-field treatment are underway[10] and an illustration is given in fig. 3 (*right*). It is seen that pions with momenta around 200 MeV/c are produced, even if the initial state contained no quasiparticles at all. The initial presence of thermal excitations will then typically increase the yield several fold.



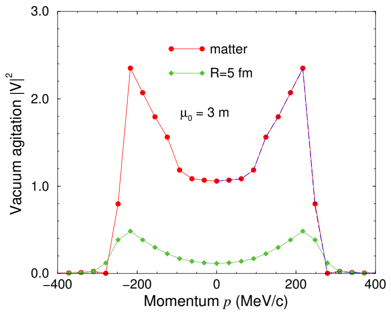


Figure 3: *Left:* The temperature dependence of the thermal occupanices of the σ and π quasiparticles as obtained with either a semi-classicial treatment[5] or an optimized perturbation method.[9] *Right:* The final pion spectrum arising solely from amplification of the vacuum fluctuations in one-dimensional systems with given effective mass functions $\mu_\pi^2(x,t)$ of forms similar to those obtained for Bjorken rods (see fig. 2 (*right*)).[10]

5 Strangeness

Since the temperatures of interest easily exceed the mass of the s quark, it is expected that strangeness must be incorporated. Within the linear σ model, this involves an extension from its familiar SU(2) form to SU(3). The number of meson fields then increases from 4 to 18,

$$\sigma;\ \pi^0, \pi^-, \pi^+ \ \longrightarrow\ \begin{cases} \sigma,\ \ a_0^0,\ \zeta; & a_0^-,\ a_0^+,\ K^{*-},\ K^{*+},\ K^{*0},\ \bar{K}^{*0}; \\ \pi^0,\ \eta_8,\ \eta_0; & \pi^-,\ \pi^+,\ K^-,\ \ K^+,\ \ K^0,\ \ \bar{K}^0 \end{cases} , \tag{4}$$

and the equilibrium value of the order parameter acquires a strange component, $\sigma \to (\sigma, \zeta)$, where $\zeta = \langle \bar{s}s \rangle$.

As is evident from Fig. 4 (*left*), the inclusion of strangeness severely impedes the restoration of chiral symmetry as the temperature is raised, thus casting doubt on the standard scenario in which approximate restoration is assumed to occur once T exceeds a few hundred MeV.

On the other hand, as illustrated in Fig. 4 (*right*), the dynamics of the order parameter becomes more intricate. As a result, the kaon modes may experience parametric amplification in analogy to what happens for the pions. The neutral kaon fraction $f_K = (K^0 + \bar{K}^0)/(K^+ + K^- + K^0 + \bar{K}^0)$ then has a uniform distribution, $P_K(f_K) = 1$, for an idealized source, while the neutral pion fraction $f_\pi = \pi^0/(\pi^+ + \pi^- + \pi^0)$ retains its SU(2) distribution.[11]

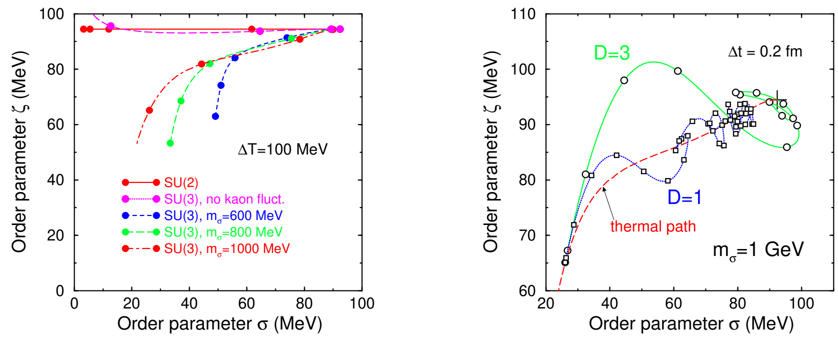


Figure 4: *Left:* The thermal path of the two order parameter (σ, ζ) for various cases of interest, as indicated. The solid dots are plotted in temperature steps of 100 MeV. All paths start from the vacuum point at the upper-right corner and the points for T=100 MeV are still very close to the vacuum point. *Right:* The dynamical path of the order parameter (σ, ζ) as a result of a pseudoexpansion in one and three dimensions, starting from equilibrium at T=400 MeV. Equidistant time steps with $\Delta\tau$=0.2 fm/c are indicated. The thermal path is shown by the long-dashed curve, while the cross indicates the vacuum point, towards which all trajectories converge in time.[11]

6 Outlook

Chiral symemtry is a fundamental concept in strong-interaction physics and hence it is of general interest to explore its properties over a wide range of environments. High-energy collisions between heavy nuclei provide a promising means for this undertaking, as dynamical explorations with the linear σ model suggest that specific signatures *may* be practically observable. It is therefore worthwhile to intensify the research efforts in this area.

On the theoretical side, many challenges remain and the present discussion has merely highlighted two of them: 1) the need for taking account of the vacuum fluctuations by developing a proper real-time non-equilibrium quantum-field treatment of the chiral dynamics and 2) the need for including the strangeness degrees of freedom into the description.

Although both problems pose considerable difficulties, formal as well as practical, the available results do provide specific guidance with respect to what may constitute particularly informative observables. These include enhancements in the soft part of the transverse pion spectra, anomalous fluctuations in their multiplicities, azimuthal correlation patterns, and a widening of the distribution of the neutral pion fraction.[7] Thus (apart from this latter observable which may be beyond early reach), there is reason to hope for the appearance of interesting data in the upcoming first round of RHIC experiments.

Acknowledgments

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References

1. V. Koch, Int. J. Mod. Phys. E6, 203 (1997).
2. K. Rajagopal, *Quark-Gluon Plasma 2* (R. Hwa), World Scientific (1995).
3. J.-P. Blaizot and A. Krzywicki, Acta Phys. Polon. B27 (1996) 1687.
4. J.D. Bjorken, Acta Phys. Polon. B28 (1997) 2773.
5. J. Randrup, *Phys. Rev.* D55, 1188 (1997); *Nucl. Phys.* A616, 531 (1997).
6. J. Randrup, *Phys. Rev. Lett.* 77, 1226 (1996).
7. T.C. Petersen and J. Randrup, *Phys. Rev.* C61, 024906 (2000).
8. J. Randrup, Heavy Ion Phys. 9, 289 (1999).
9. S. Chiku and T. Hatsuda, *Phys. Rev.* D58, 076001 (1998).
10. J. Randrup, LBNL-00000 (2000) in preparation.
11. J. Schaffner-Bielich and J. Randrup, *Phys. Rev.* C59, 3329 (1999).