Glauber dynamics of phase transitions: SU(3) lattice gauge theory.

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Motivated by questions about the QCD deconfining phase transition, we studied in two previous papers Model A (Glauber) dynamics of 2D and 3D Potts models, focusing on structure factor evolution under heating (heating in the gauge theory notation, i.e., cooling of the spin systems). In the present paper we set for 3D Potts models (Ising and 3-state) the scale of the dynamical effects by comparing to equilibrium results at first and second order phase transition temperatures, obtained by re-weighting from a multicanonical ensemble. Our finding is that the dynamics entirely overwhelms the critical and non-critical equilibrium effects.

In the second half of the paper we extend our results by investigating the Glauber dynamics of pure SU(3) lattice gauge on $N_x N_y N_z$ lattices directly under heating quenches from the confined into the deconfined regime. The exponential growth factors of the initial response are calculated, which give Debye screening mass estimates. The quench leads to competing vacuum domains of distinct $Z_3$ triality, which delay equilibration of pure gauge theory forever, while their role in full QCD remains a subtle question. As in spin systems we find for pure SU(3) gauge theory a dynamical growth of structure factors, reaching maxima which scale approximately with the volume of the system, before settling down to equilibrium. Their influence on various observables is studied and different lattice sizes are simulated to illustrate an approach to a finite volume continuum limit. Strong correlations are found during the dynamical process, but not in the deconfined phase at equilibrium.

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I. INTRODUCTION

In investigations of the QCD deconfining phase transition (or crossover) by means of heavy ion experiments, one ought to be concerned about non-equilibrium effects due to the rapid heating of the system. With this in mind we have investigated in previous papers\textsuperscript{1,2} the Model A\textsuperscript{3} (Glauber) dynamics of 2D and 3D Potts models. Model A dynamics includes all diffusive stochastic local updating schemes (Metropolis, heatbath, etc.) and not only the process introduced in\textsuperscript{4}. In 3D Potts models spins provide degrees of freedom, which mimic Polyakov loops effectively\textsuperscript{5}, while in 2D analytical results\textsuperscript{6} allow to check on the accuracy of the employed numerical methods. For other approaches to simulate non-equilibrium quantum fields see Ref.\textsuperscript{7}.

The QCD high temperature vacuum is characterized by ordered Polyakov loops, which are similar to spins in the low temperature phase of the 3D 3-state Potts model. We model heating by a quench from the disordered into the ordered phase, which thus corresponds to a cooling quench in the analogue spin model. Time evolution after the quench leads to vacuum domains of distinct triality under the $Z_3$ center of the SU(3) gauge group. It appears that these competing domains are the underlying cause for the explosive growth of structure factors $F_i(t)$, which we encounter in the time evolution after a heating quench. We use the term spinodal decomposition loosely to denote generically such a time period of globally unstable behavior.

Relaxation of the system at its new temperature becomes only feasible after each structure factor has over-
3D Ising and 3-state Potts model. Subsequently their dynamical evolution after a quench is investigated, extending previous results. In section IV we present our simulations of pure SU(3) lattice gauge theory. Some SU(3) data were already reported at the 2004 APS DPF conference 10. As these simulations are very CPU time consuming it took over one more year to collect the present statistics. Summary and conclusions are given in the final section V.

II. NOTATION AND PRELIMINARIES

We summarize our basic notations and concepts in this section.

A. Models

We simulate q-state Potts models with the energy function

\[ E = 2 \sum_{\langle ij \rangle} \left( \frac{1}{q} - \delta_{q_i q_j} \right) \]  

(1)

where the sum is over nearest neighbors of a hypercubic lattice in D dimensions. The spins \( q_i \) of the system take on the values \( q_i = 0, \ldots, q - 1 \). The factor of two and the term \( 1/q \) is introduced to match for \( q = 2 \) with Ising model conventions 11. Simulations are carried out with the Boltzmann factor \( \exp(-\beta E) \).

The Wilson action for pure SU(3) non-Abelian Euclidean lattice gauge theory is

\[ S_A = \frac{2 \cdot 3}{g^2} \sum_{n, \mu \nu} \left[ 1 - \frac{1}{2} \cdot \text{Tr}(U_{n, \mu \nu} + \text{h.c.}) \right] \]  

(2)

where \( U_{n, \mu \nu} = U_{n, \mu} U_{n+\hat{\mu}, \nu} U_{n+\hat{\nu}, \mu} U_{n+\hat{\mu}+\hat{\nu}, \nu} \) denotes the product of the SU(3) link matrices in the fundamental representation around a plaquette and the sum runs over all plaquettes. Simulations are done with the Boltzmann factor exp(\( S_A \)).

The Markov chain Monte Carlo (MC) process provides model A (Glauber) dynamics in the classification of Ref. 3. For Potts models we use the heatbath algorithm of 12 and for SU(3) gauge theory the Cabibbo-Marinari 13 heatbath algorithm and its improvements of Ref. 13 (no over-relaxation, to stay in the universality class of Glauber dynamics). In each case a time step is a sweep of systematic updating through the lattice, which touches each degree of freedom once. With small statistics we have checked that updating in random order gives similar results up to a slowing down of the evolution speed by a constant factor. This is expected as in equilibrium simulations random updating has larger autocorrelations than systematic updating 11. For our equilibrium simulations of Potts models we used a multicanonical 14 Metropolis algorithm.

B. Structure Factors

Consider two-point correlation functions defined by

\[ \langle u_0(0) u^\dagger_0(j) \rangle_L = \frac{1}{N^2} \sum_i u_0(i) u^\dagger_0(i + j), \]  

(3)

where \( \vec{i} \) denotes spatial coordinates. Periodic boundary conditions are used and the subscript \( L \) on the left-hand side reminds us that the average is taken over the spatial lattice. For gauge systems we deal with fluctuations of the Polyakov loop, for analogue spin systems with fluctuations of the magnetization.

The finite volume continuum limit of (3) is achieved by lattice spacing \( a \to 0 \) and \( N_L \to \infty \) with the physical length of the box \( L = a N_L \) = const. This means that

\[ \langle u_0(0) u^\dagger_0(j) \rangle_L = \frac{1}{a^2 N_L} \sum_i a^3 u_0(i) u^\dagger_0(i + j) \]  

(4)

transforms into

\[ \langle u(0) u^\dagger(\vec{R}) \rangle_L = \frac{1}{L^3} \int d^3 \vec{r} u(\vec{r}) u^\dagger(\vec{r} + \vec{R}), \]  

(5)

with \( \vec{r} = a \vec{i}, \vec{R} = a \vec{j}, u(\vec{r}) = u_0(\vec{i}), \) and so on. We define the structure function \( F(\vec{p}) \) as Fourier transformation of the two-point correlation function (5):

\[ F(\vec{p}) = \int \langle u(0) u^\dagger(\vec{R}) \rangle_L e^{i \vec{p} \cdot \vec{R}} d^3 \vec{R}. \]  

(6)

Periodic boundary conditions imply:

\[ \vec{p} = \frac{\vec{k}}{a} = \frac{2 \pi}{L} \vec{n}, \]  

(7)

where \( \vec{n} \) is an integer vector \((0, 0, 0), (0, 0, 1), \) and so on. The discretized version of (6) is

\[ F(\vec{p}) = \sum_j a^3 \langle u_0(0) u^\dagger_0(j) \rangle_L e^{i \vec{p} \cdot \vec{j}}. \]  

(8)

Using the definition (8) and shifting the \( \vec{j} \) summation one arrives (after straightforward algebra) at the expression

\[ F(\vec{p}) = \frac{a^3}{N^2} \left| \sum_i e^{-i \vec{p} \cdot \vec{i}} u_0(\vec{i}) \right|^2, \]  

(9)

where we may rewrite the product in the exponent as

\[ a \vec{p} \vec{i} = \vec{k} \vec{i} = \frac{2 \pi}{N} \vec{n} \vec{i}. \]  

(10)

As we let the system evolve after a quench \( u_0(\vec{i}) \) becomes time-dependent: \( u_0(\vec{i}, t) \). The time \( t \) corresponds to the dynamical process, i.e., in our case the Markov chain model A dynamics. We consider an ensemble of systems
(replica) and dynamical observables are calculated as ensemble averages denoted by \( \langle ... \rangle \). The time-dependent structure functions averaged over replicas are:

\[
F_p(t) = \langle F(\tilde{p}, t) \rangle .
\]  

During our simulations they are averaged over rotationally equivalent momenta and the notation

\[
F_i(t)
\]

is used for the structure function at momentum

\[
\tilde{p} = \tilde{k} = \frac{2\pi}{L} \tilde{n}
\]

where \(|\tilde{n}| = n_i \) defines \( i \). The \( F_i \) are called structure function modes or structure factors (SFs). We recorded the following modes (including the permutations) \( n_1: (1, 0, 0), n_2: (1, 1, 0), n_3: (1, 1, 1), n_4: (2, 0, 0), n_5: (2, 1, 0), n_6: (2, 1, 1), n_7: (2, 2, 0), n_8: (2, 2, 1) \) and \( (3, 0, 0), n_9: (3, 1, 0) \). Note that there is an accidental degeneracy in length for \( n_8 \). We measured also higher modes in some cases up to \( n_{64} \). They exhibit the same behavior as the lower modes, but the data are far more noisy, so that we abstain from reporting these results. A difference to the normalization of \( \| \tilde{x} \| \) is that in the present paper we average over the permuted momenta instead of just summing them up. For instance, for the \( F_1 \) SF the difference is a multiplicative factor of three.

### III. POTTS MODELS

For the analogue spin models the lattice spacing \( a \) cannot be varied. We set \( a = 1 \), so that the distinction between \( L \) and \( N_s \), \( \tilde{p} \) and \( \tilde{k} \) becomes superfluous. We use \( L \) and \( \tilde{k} \) in the following. The normalization of the SFs differs from our previous work [1, 2]. It is chosen so that they approach constant values in the infinite volume limit of equilibrium simulations of spin systems at non-critical temperatures. This follows from the fact that the random fluctuations in \( \| \tilde{x} \| \) are of order \( \sqrt{V} = \sqrt{L^3} \). At a critical temperature of a second order phase transition a divergence of the SFs is then encountered as we illustrate for the 3D Ising model. A sustained increase of a SF with lattice size cannot be stronger than being proportional to \( V = L^3 \), because an upper bound on each SF is obtained by setting all values in the sum of Eq. (10) equal to one.

#### A. Equilibrium Results

In this section we compile SF estimates from equilibrium simulations of the 3D Ising and 3-state Potts model on \( L^3 \) lattices. Our simulations are carried out in a multicanonical ensemble [13], covering a temperature range from \( \beta_{\text{min}} = 0 \) (infinite temperature) to \( \beta_{\text{max}} > 0 \) below the phase transition temperature of the respective system. Instead of relying on a recursion (see, e.g., [11]), the multicanonical parameters were extracted by finite size (FS) extrapolation from smaller to larger system, which is an efficient way when the FS behavior is controllable.

The advantage of using multicanonical simulations is that accurate values of the SF peaks can be determined from one data set. Re-weighting of a canonical simulation [12] allows accurate determination of the maxima of one quantity, but on finite lattices the maxima of different observables are too far apart to be within the re-weighting range of one canonical simulation. We find it convenient to have the entire range of interest covered in one simulation. In particular equilibration of the configurations around the transition and in the ordered phase is then secured due to frequent excursions into the discarded region all the way to \( \beta = 0 \).

#### 1. 3D Ising Model

At the critical point the two-point function on an infinite lattice falls off with a power law, which defines the critical exponent \( \eta \):

\[
f(\tilde{x}) = \langle s(\tilde{x}) s(\tilde{x}') \rangle \sim |\tilde{x}|^{-d+2-\eta}, \quad |\tilde{x}| \to \infty.
\]  

This determines the low-momentum behavior of the Fourier transformation \( F(\tilde{k}) \). Namely,

\[
\hat{F}(\lambda \tilde{k}) \sim \int d^d x e^{i \lambda \tilde{k} \cdot \tilde{x}} |\tilde{x}|^{-d+2-\eta} = \int \frac{d^d x'}{\lambda^d} e^{i \tilde{k} \cdot \tilde{x}'} \lambda^{d-2+\eta} |\tilde{x}'|^{-d+2-\eta} \sim \lambda^{\eta-2} \hat{F}(\tilde{k})
\]

holds and, therefore,

\[
\hat{F}(\tilde{k}) \sim |\tilde{k}|^{-b}, \quad b = 2 - \eta, \quad \text{for} \quad |\tilde{k}| \to 0.
\]  

For fixed \( \tilde{n} \) we have \( \tilde{k} = 2\pi \tilde{n}/L \) and we find for any fixed value of \( \tilde{n} \) the finite size scaling (FSS) divergence

\[
\hat{F}_{\tilde{n}} \sim L^b, \quad b = 2 - \eta, \quad \text{for} \quad L \to \infty
\]  

with lattice size.

---

**TABLE I: Statistics and SF maxima \( F_{i,\text{max}} \) at \( \beta_m \) from our equilibrium simulations of the 3D Ising model on \( L^3 \) lattices.**

| \( L \) sweeps | \( F_{i,\text{max}}^{\text{max}} \) | \( \beta_m \) | cycles |
|----------------|-----------------|--------|-------|
| 20 | \( 32 \times 5 \times 10^4 \) | 17.00 (26) | 0.219874 | 83 |
| 30 | \( 32 \times 2 \times 10^5 \) | 37.97 (70) | 0.220825 | 89 |
| 44 | \( 32 \times 6 \times 10^5 \) | 78.4 (1.6) | 0.221146 | 70 |
| 56 | \( 32 \times 1 \times 10^6 \) | 129.6 (2.7) | 0.221345 | 45 |
| 66 | \( 32 \times 1.6 \times 10^6 \) | 175.6 (4.9) | 0.221387 | 43 |
| 80 | \( 64 \times 2 \times 10^6 \times 3 \) | 257.4 (2.4) | 0.221462 | 65 + 72 + 67 |
The infinite volume phase transition temperature of the 3D Ising model is estimated to be $\beta_c = 0.22157 (3)$, e.g., see [16]. In our multicanonical simulations we cover the range from $\beta_{\text{min}} = 0$ to $\beta_{\text{max}} = 0.25$, well including the transition region. Table I gives an overview of the lattice sizes and the accumulated statistics (a sweep updates each spin once) together with our estimate of the maximum values $F_1^{\text{max}}$ of the first SF, evaluated at the value $\beta_m$. Error bars are given in parenthesis and apply to the last digits of the number in front. They are calculated with respect to a number of jackknife bins given by the first number in column two of the table, and the multicanonical re-weighting procedure uses the logarithmic coding described in [17]. Three independent runs were carried out for the $L = 80$ lattice. Before starting with measurements we normally performed the number of sweeps of one measurement bin for reaching equilibrium. This is sufficient because equilibration problems are mild in multicanonical simulations. Running time for each of our $L = 80$ simulations was about three months on a 2 GHz Athlon PC. The last column of table I gives the number of cycles $(\beta_e \leq \beta_{\text{min}}) \rightarrow (\beta_e \geq \beta_{\text{max}}) \rightarrow (\beta_e \leq \beta_{\text{min}})$, which the Markov process performed during the production run, where $\beta_e$ is the effective energy-dependent $\beta$ of the multicanonical procedure.

For our largest volume the SFs 1-9 are plotted in Fig. 1, where we restrict $\beta$ to a neighborhood of the critical temperature. Each SF develops a clear peak, only that the peaks for the higher SFs are less pronounced than those for the lower. In particular the scale of the figure does not resolve the peaks for the SFs $\geq 7$ anymore. These peaks are found on a reduced scale and for each SF the FSS behavior (16) holds. However, the numerical accuracy decreases with increasing $|\mathbf{k}|$. So we are content with simply analyzing the FSS behavior of SF 1. Fig. 2 shows SF 1 for all our lattice sizes and the maxima values are collected in table I. A two parameter fit to the form (16) is shown in Fig. 3. It gives $b = 1.959 (12)$ with a goodness of fit $Q = 0.82$ (for the definition of $Q$ see, e.g., Ref. [11]), a result well compatible with the high precision estimates $\eta = 0.0364 (5)$ given in the review article [15] on critical phenomena and renormalization group theory.

2. 3D 3-state Potts Model

For the 3D 3-state Potts model one deals with a relatively weak first order phase transition at $\beta_c = 0.2752720 (49)$, a value which averages two somewhat inconsistent ($Q = 0.003$ for the Gaussian difference test) estimates of the literature [10] (because of the inconsistency the error bars are averaged here and not reduced). In our multicanonical simulations we cover the range from
An overview of the statistics and some results are given in table II similarly as before for the 3D Ising model in table I.

For our $50^3$ lattice the SFs 1-9 are plotted in Fig. 4, where we restrict $\beta$ to a neighborhood of the transition temperature. As for the 3D Ising model each SF develops a clear peak, but the shapes are significantly different. A relatively smooth increase is followed by a rather abrupt decrease. The lattice size dependence of SF 1 is depicted in Fig. 5 which indicates (as expected) that the abrupt decrease develops into a discontinuity for $L \to \infty$. The increase of the structure function maxima is irregular and smaller from $L = 40$ to $L = 50$ than from $L = 30$ to $L = 40$. Asymptotically for $L \to \infty$ a finite maximum value is expected in case of a first order phase transition.

Within our limited lattice sizes this is not yet seen, but a power law fit of the type of Fig. 3 which is the large $L$ behavior in case of a second order transition, becomes entirely inconsistent: $Q = 2.7 \cdot 10^{-11}$ is the goodness of fit obtained.

### B. Quenches

After outlining the equilibrium scenario, let us discuss the time evolution after a quench from the disordered into the ordered phase of the 3D 3-state Potts model.

### FIG. 4: SFs $F_i$, $i = 1, \ldots, 9$ from simulations of the 3-state Potts model on a $50^3$ lattice.

### FIG. 5: Finite size behavior of SF $F_1$ from 3-state Potts model simulations on $L^3$ lattices.

### TABLE II: Statistics and SF maxima $F_1^{\text{max}}$ at $\beta_m$ for our equilibrium simulations of the 3D 3-state Potts model on $L^3$ lattices.

| $L$ | sweeps | $F_1^{\text{max}}$ | $\beta_m$ | cycles |
|-----|--------|---------------------|----------|--------|
| 20  | $32 \times 1.2 \cdot 10^2$ | 19.00 (22) | 0.274273 | 59     |
| 30  | $32 \times 5.2 \cdot 10^2$ | 38.11 (41) | 0.274924 | 71     |
| 40  | $32 \times 1.5 \cdot 10^2$ | 60.30 (50) | 0.275116 | 73     |
| 50  | $126 \times 1.5 \cdot 10^2$ | 80.46 (55) | 0.275181 | 131    |

### TABLE III: Repetitions of quenches from $\beta = 0.2$ to $\beta_f$ for the 3D 3-state Potts model on $L^3$ lattices.

| $\beta_f \backslash L$ | 20 | 40 | 60 | 80 | 100 | 120 |
|----------------------|----|----|----|----|-----|-----|
| 0.3                  | 640| 640| 640| 320| 320 | 320 |
| 0.27                 | 640| 320| 320| 32 |   |    |

An overview of our statistics is given in table III. We quench from $\beta = 0.2$ to the $\beta_f$ value given in the table, which collects the numbers of repetitions of each quench. Error bars are then calculated with respect to 32 bins. Larger lattices exhibit self-averaging, so that one needs less repetitions than for smaller lattices.

In previous work [2] we have investigated the quench $\beta = 0.2 \to 0.3$ and its subsequent stochastic time evolution on lattices up to size $80^3$. Meanwhile we have extended the SF part of this investigation to lattices of size up to $120^3$ and Fig. 6 shows the time evolution of SF 1 after this quench. Note that we divide the SF by an extra volume factor in this figure. So its initial increase may be expected: The quench changes the temperature in the entire lattice instantaneously. It is then plausible that the local contribution to the SF is, in the average, everywhere the same. So one expects an increase $\sim V$ of the maxima. The initial overshooting may be explained by an increase of correlations with lattice size, which levels off once the lattice size exceeds the correlation length.

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Heuristically this behavior during spinodal decomposition may be expected: The quench changes the temperature in the entire lattice instantaneously. It is then plausible that the local contribution to the SF is, in the average, everywhere the same. So one expects an increase $\sim V$ of the maxima. The initial overshooting may be explained by an increase of correlations with lattice size, which levels off once the lattice size exceeds the correlation length.

To test how this growth of the signal proportional to the volume depends on the depth of the quench into the
ordered region, we performed a quench to a temperature closer to the transition temperature, $\beta = 0.2 \to 0.28$. As shown in Fig. 7 we find the same phenomenon as before: The maximum sustained increase $\sim V$ is initially over-shot. The growth of the signal is weaker than before, as is expected since the system does not change so drastically.

In both figures we see that the time positions $t_{\text{max}}$ of the SF $F_1$ maxima move towards larger values with increasing lattice size. For our two quenches $t_{\text{max}}(L)$ is plotted in Fig. 8 on a log-log scale. With parameters $a_0$ and $a_1$ both curves can consistently be fitted to the expected form

$$t_{\text{max}}(L) = a_0 + a_1 L^2.$$  \hfill (18)

As $t$ is measured in units of sweeps, the number of spin updates per time unit does not depend on $L$. In spin systems $t$ is thus proportional to the physical time. After the quench into the ordered phase the infinite spin system

cannot be equilibrated in any finite time, a fact known in condensed matter physics [20]. The explanation for this phenomenon is that the systems grows initially competing domains of three distinct orientations. To dissolve these domains by local random fluctuations until one of them dominates the entire lattice is a slow process, which requires of order $L^2$ time.

Visualization of these domains faces difficulties, because naive geometrical definitions do not work. Compare Fig. 8 of Ref. [2]. For analogue Potts models this is overcome by the Fortuin-Kasteleyn [21] cluster definition, but there is no immediate generalization to gauge theories, although promising ideas have been published [22]. Here we do not investigate this question any further. We think that it is safe to assume that competing domains are in both, spin and gauge systems, the underlying cause for the explosive growth of structure factors $F_i(t)$, which we encounter in their time evolution after a heating quench.

Finally in this section, based on 640 repetitions Fig. 9 demonstrates that for a non-critical quench nothing more than a smooth transition from one equilibrium value to the next happens. Therefore the explosive growth of SFs is a unambiguous signal that $\beta_f$ is indeed in the ordered phase.

IV. SU(3)

We report results from quenches of pure SU(3) lattice gauge theory on $N_{\tau} N_{\sigma}^3$ lattices. Our statistics is summarized in tables IV and V. All quenches are from the initial value $6/g^2 = 5.5$. The $4 \times N_{\sigma}^3$ simulations of table IV were already reported in [14]. The simulations for the other lattices are new. The difference between the tables is that for the lattices of table IV we follow the quench all the way to its equilibrium value at $T_f$, while for the lattices of table V we calculated only the initial
increase of the SFs as needed for the determinations of critical modes in section IV.B.

The new data serve to study the quantum continuum limit $a \to 0$ (in physical units like fermi). The final values $g_f^2$ of the bare coupling constants are chosen, so that the values of $T_f/T_c$ stay at the fixed ratios given in the table. For this we take (substantial) corrections to the two-loop equation of Lambda lattice into account, which follow from renormalization group results tabulated in Ref. [28]. As the use of tables is tedious, we like to mention that with an accuracy of 0.5% and better our $T_f/T_c$ values are reproduced by using the formula

$$\Lambda_L(g^2) = \Lambda_L^{as}(g^2) \lambda(g^2)$$  \hspace{1cm} (19)

where $\Lambda_L^{as}(g^2)$ is given by (e.g., [24])

$$\Lambda_L^{as} = (b_0 g_0^2)^{-b_1/(2b_0)} e^{-1/(2b_0 g^2)}$$

with $b_0 = \frac{11}{3} \frac{N}{16\pi}$, $b_1 = \frac{33}{3} \left(\frac{N}{16\pi}\right)^2$ and

$$\lambda(g^2) = 1 + a_1 e^{-a_2/g^2} + a_3 g^2 + a_4 g^4$$

with $a_1 = 71553750$, $a_2 = 19.48099$, $a_3 = -0.03772473$, and $a_4 = 0.5089052$.

For $N_t = 4$, fixed, Fig. 10 shows the divergence of the SF 1 maxima with increasing lattice size $N_v^3$ as well as the $f_{\text{max}}(N_v) \sim N_v^2$ behavior in complete analogy to our results for Potts models. All the lattices of Fig. 10 are quenched to the bare coupling constant $g_f^2 = 6/5.92$. Therefore the time scale of the Markov process (determined by the Boltzmann factors) is the same on all these lattices and up to an unknown multiplicative factor identified with that of a dissipative, non-relativistic dynamics. Non-relativistic does not necessarily mean that the propagation of the signal through the lattice is slow. In the contrary, Galilee transformations set no upper limit on speeds. Our quench changes the temperature instantaneously through the entire lattice, while the subsequent propagation of the response proceeds through local interactions.

A. Finite Volume Continuum Limit

In the following we illustrate the approach of the limit $a \to 0$, $L = \text{constant}$, $T_f/T_c = \text{constant}$, by increasing $N_v$ from 4 to 6 to 8 and the volume $N_v^3$ from $N_v = 16$ to 24 to 32, so that the ratio $N_v/T$ stays constant. The ratio of temperatures $T_f/T_c$ is kept constant by using Eq. (19) to determine the appropriate bare coupling
bare coupling constants values for different situation changes too, because we have to use different
the time axis to finite time, we fix this normalization problem by rescaling

$N_\tau$ lattices of constant physical volume of a
quench to $T_f/T_c = 1.25$. of our original measurements in units of sweeps are related to those used in Fig. 11 by the $\lambda_i(N_\tau,1.25)$ factors 1 : 2.655 : 5.457 for $N_\tau$ the values 4 : 6 : 8, respectively. For Fig. 12 the corresponding $\lambda_i(N_\tau,1.568)$ ratios are 1 : 2.768 : 6.362. The maxima of the curves decrease when increasing $N_\tau$ from 4 to 6 to 8. As the decrease slows down with increasing lattice size, there is some evidence for an approach to a shape, which represents the continuum limit.

B. Debye Screening Mass

The current understanding of the early time evolution of systems out of equilibrium is largely based on investigating stochastic equations which are dynamical (time dependent) generalizations of the Landau-Ginzburg effective action models of the static (equilibrium) theory. For model A the linear approximation results in the following equation for a SF:

$$\frac{\partial \hat{F}(\vec{p},t)}{\partial t} = 2 \omega(\vec{p}) \hat{F}(\vec{p},t),$$

with the solution

$$\hat{F}(\vec{p},t) = \hat{F}(\vec{p},t=0) \exp(2\omega(\vec{p})t),$$

where $p_c > 0$ is a critical momentum. Originally the linear theory was developed for model B. Details for model A can be found in Ref. [25].

From our measurements of $F(\vec{p},t)$ on the $N_\tau = 4$, 6 and 8 lattices we find straight line fits to the form

$$\omega(p) = a_0 + a_1 p^2, \quad p = |\vec{p}|,$$ with a negative slope $a_1$. They determine the critical momentum $p_c$ as the value where

$\omega(p)$ changes its sign. The fits for $T_f/T_c = 1.25$ are shown in Fig. 13 and for $T_f/T_c = 1.568$ in Fig. 14, where

FIG. 11: Time evolution of SF $F_i/F_{i,f}$ for SU(3) lattice gauge theory on $N_\tau N_\sigma^3$ lattices of constant physical volume of a quench to $T_f/T_c = 1.568$.

FIG. 12: Time evolution of SF $F_i/F_{i,f}$ for SU(3) lattice gauge theory on $N_\tau N_\sigma^3$ lattices of constant physical volume of a quench to $T_f/T_c = 1.568$.

FIG. 13: SU(3) determination of $p_c$ for $T_f/T_c = 1.25$. we overcome by dividing all SFs $F_i$
by their equilibrium values at $T_f$, $F_{i,f}$. The time-scale situation changes too, because we have to use different bare coupling constants values for different $N_\tau$. As one knows that a finite physical volume equilibrates in a finite time, we fix this normalization problem by rescaling the time axis to

$$t' = \frac{t}{\lambda_i(N_\tau,T_f/T_c)}$$ (20)

so that all maxima fall on top of one another. We do not lose information as we anyhow do not know the overall normalization factor for our time scale.

FIG. 14: Determination of $\frac{T_c}{T_f}$ at $N_\tau = 4$, 6 and 8 lattices.
we introduced
\[ \omega(p) = \lambda_t(N_r, T_f/T_c) \omega(p) \]  
(23)
This definition absorbs the shift \( \omega'(p) t' = \omega(p) t \) holds. It is only in the primed variables that one realizes an eventual approach to the continuous limit from Figs. 13 and 14. In particular note that for \( T_f/T_c = 1.568 \) the \( N_r = 6 \) and 8 fits are within statistical errors identical. The obtained values for \( p_c(N_r)/T_c \) are listed in table VI. The (finite volume) continuum limit is extrapolated by fitting these values to the form
\[ \frac{p_c(N_r)}{T_c} = \frac{p_c}{T_c} + \frac{\text{const}}{N_r} \]  
(24)
with the results given in the last column of table VI.

Relying on a phenomenological analysis by Miller and Ogilvie, \( p_c \), is related by
\[ m_D = \sqrt{3} p_c \]  
(25)
to the Debye screening mass at the final temperature \( T_f \) after the quench. We get
\[ m_D = 1.77 (15) T_c \text{ for } T_f/T_c = 1.25, \]  
(26)
\[ m_D = 3.53 (13) T_c \text{ for } T_f/T_c = 1.568. \]  
(27)
The value at \( T_f/T_c = 1.568 \) is in excellent agreement with a determination of \( m_D(T) \) from a best-fit analysis of the large distance part of the color singlet free energies. \( p_c \) in two smaller than the one of Ref. 29. This is not really a surprise, because \( T_f/T_c = 1.25 \) is close to the spinodal endpoint, so that the derivation of the relationship 29 is no longer valid.

For pure SU(3) lattice gauge theory \( T_c = 265 \) (1) MeV holds, assuming \( \sigma = 420 \) MeV for the string tension, while for QCD the cross-over temperature appears to be around \( T_c \approx 165 \) MeV, see Ref. 30 for reviews. Using for simplicity \( T_c = 200 \) MeV to illustrate the magnitudes, the temporal lattice size is then about 1 fermi at \( T_c \). The spatial sizes of our lattices used in this section reach up to \( (8 \text{ fermi})^3 \). At the \( T_f \) values the edge lengths are shortened by the corresponding \( T_c/T_f \) factors. I.e., the volume is \( (6.4 \text{ fermi})^3 \) for \( T_f/T_c = 1.25 \) and \( (5.10 \text{ fermi})^3 \) for \( T_f/T_c = 1.568 \). The screening length associated with the Debye mass, \( \xi_D = 1/m_D, \) is then approximately 0.6 fermi at \( T_f/T_c = 1.25 \) and 0.3 fermi at \( T_f/T_c = 1.568 \). The illustration of the finite volume continuum limit in section VI was for lattices of size \( (4 \text{ fermi})^3 \) at \( T_c \), i.e., \( (3.2 \text{ fermi})^3 \) at \( T_f/T_c = 1.25 \) and \( (2.55 \text{ fermi})^3 \) at \( T_f/T_c = 1.568 \). Our volumes are smaller than the envisioned deconfined region of about \( (10 \text{ fermi})^3 \) in relativistic heavy ion experiments. Due to periodic boundary conditions one may expect that MC simulations on smaller lattices are representative for the central region of the larger volume. Our result is that the Debye screening length is short on the scale of the deconfined region.

C. Measurements near Structure Factor Maxima versus Deconfined Equilibrium

For SU(3) gauge theory the triality of Polyakov loops with respect to the \( Z_3 \) center of the gauge group takes the place of three distinct spin orientation. Although a satisfactory cluster definition does not exist for gauge theories, the underlying mechanism of competing vacuum domains is expected to be similar as in the spin models. To study their influence on Polyakov loop correlations and on the gluonic energy \( \epsilon \) and pressure \( p \) densities, we calculate these quantities at times \( t \leq t_{\text{max}} \) as well as at \( t > t_{\text{max}} \).

Our structure function measurements gave “on the fly” two-point correlations between Polyakov loops defined by
\[ C_o(d,t) = \langle P(0,t) P(d,t) \rangle_L = \langle |\langle P(0,t)\rangle_L|^2 \rangle \]  
(28)
where the averaging procedures are those we discussed after Eq. 13 and \( d = 1, 2, \sqrt{2}, \ldots \). The value of these results is somewhat limited, because our focus was not on good equilibrium results and the stored data do not allow to project onto particular channels of the free energy of static quarks (which lead to larger correlation lengths than those obtained). For several values of \( d \) we plot in Fig. 14 the time development of \( C_o(d) \) on our largest lattice using the \( t' \) time scale. The correlations assume maxima at about the same time values \( t_{\text{max}} \) for which the SFs peak, although less pronounced.
In Fig. 15 we plot the $d$-dependence for the time values $0.5t_{\text{max}}$, $t_{\text{max}}$ and $5t_{\text{max}}$. At $5t_{\text{max}}$ fits of the form $C_o(d) \sim \exp(-m_P a d)/(a d)$, where $a$ is the lattice spacing, give the $m_P$ estimates which are collected in table VII. The last column of this table gives infinite volume estimates obtained from fits of the form (24). In contrast to that large correlations are found at $0.5t_{\text{max}}$ and $t_{\text{max}}$, which are fully consistent with a power law.

The equilibrium procedure for calculating the gluonic energy $\epsilon$ and pressure $p$ densities is summarized in Ref. 23, 31 (in earlier work 32 the pressure exhibited a non-physical behavior after the deconfining transition and the energy density approached the ideal gas limit too quickly because the anisotropy coefficients were calculated perturbatively). We denote expectation values of space-like plaquettes by $P_\sigma$ and those involving one time link by $P_\tau$. The energy density and pressure can then be cast into the form
\[
\frac{\epsilon + p}{T^4} = \frac{8 N_c N_T^4}{g^2} \left[ 1 - \frac{g^2}{2} [c_\sigma(a) - c_\tau(a)] \right] (P_\sigma - P_\tau)
\]
and
\[
\frac{\epsilon - 3p}{T^4} = 12 N_c N_T^4 \left[ c_\sigma(a) - c_\tau(a) \right] \left[ 2 P_0 - (P_\sigma + P_\tau) \right],
\]
where $P_0$ is the plaquette expectation value on a symmetric $(T = 0)$ lattice, and the anisotropy coefficients $c_{\sigma,\tau}(a)$ are defined by:
\[
c_{\sigma,\tau}(a) \equiv \left( \frac{\partial g_\sigma^{-2}}{\partial \xi} \right)_{\xi=1}.
\]

They are related to the QCD $\beta$-function and can be calculated using Padé fits of 22. To normalize to zero temperature, plaquette values from the symmetric $N_r = N_c$ lattice are needed in Eq. (30). As one stays within the confined phase on the symmetric lattice its equilibration after the quench is fast. Therefore it is enough to use equilibrium values of $P_0$ at $\beta_f$ after the quench. This is illustrated in Fig. 17.

In Fig. 18 we show the time evolution of the gluonic energy densities (upper curves) and pressure densities (lower curves) for the $T_f/T_c = 1.25$ quench on our $4 \times 16^3$, $6 \times 24^3$ and $8 \times 32^3$ lattices using the rescaled time definition 20. The curves for the last two lattices fall almost...
V. SUMMARY AND CONCLUSIONS

In equilibrium at temperatures much higher than the deconfinement temperature $T_c$ the perturbative prescription of QCD is that of a weakly coupled gas of quasiparticles. In contrast to that recent experiments at the BNL relativistic heavy-ion collider (RHIC) show coherence in particle production and strong collective phenomena, which are well described by the model of a near-perfect, strongly coupled fluid [33]. Non-perturbative effects are expected to play some role in the prescription of equilibrium QCD at temperatures reached at the RHIC. For the $T_f/T_c = 1.25$ and $T_f/T_c = 1.568$ temperatures investigated in this paper equilibrium lattice calculations indicate indeed corrections (compare Fig. 7 of [23]). However, correlations are typically over ranges much smaller than the deconfined region, compare our estimates of the Debye screening mass $m_D(T_f)$. The agreement of our $m_D$ value at $T_f/T_c = 1.568$ with direct equilibrium estimates [29] give confidence that model A dynamics reflects physical features.

If the phenomenological description of a strongly coupled plasma implies correlations over distances exceeding one fermi, the time evolution of our structure factors (SFs) depicted in Fig. 6 and 11 suggest a scenario in which the deconfined equilibrium phase has actually not been reached at the RHIC, but the heating process gets stuck during the time period of explosive growth of the SFs. While this explains correlation over distances much larger than one fermi, it also provides an unambiguous signal for the existence of the deconfining phase: Fig. 10 demonstrates that the explosive growth is absent for a non-critical quench.

In real QCD there are two effects which prevent the divergence of the equilibration time shown in Fig. 5: (1) Quarks break the $Z_3$ symmetry of the SU(3) gauge group, similarly as a magnetic field breaks the degeneracy of the spins in the 3D 3-state Potts model. The final magnitude of the equilibration time depends then on the strength of the breaking as illustrated in Ref. 2 for a weak magnetic field. (2) At the RHIC the physical volume is finite, so that even in case of an exact symmetry the equilibration time is finite. So the scenario that the system gets stuck during the spinodal decomposition of its vacuum structure could only be based on phenomenological observations. Questions like how a perfect fluid may look during the period of spinodal decomposition arise. Minkowski space simulations of hyperbolic differential equations, which emerge from effective actions for Polyakov loops [34, 35], may shed light on the question whether features observed in the paper are special to Glauber dynamics or of some universal nature.

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[1] B.A. Berg, U.M. Heller, H. Meyer-Ortmanns, and A. Velytsky, Phys. Rev. D 69, 034501 (2004).
[2] B.A. Berg, H. Meyer-Ortmanns, and A. Velytsky, Phys. Rev. D 70, 054505 (2004).
[3] P.M. Chaikin and T.C. Lubensky, Principles of condensed matter physics (Cambridge University Press, Cambridge, 1997), table 8.6.1, p.467.
[4] R.J. Glauber, J. Math. Phys. 4, 294 (1963).
[5] B. Svetitsky and L.G. Yaffe, Nucl. Phys. B 210, 443 (1982).
[6] R.J. Baxter, J. Phys. C 8, L445 (1973); F.Y. Wu, Rev. Mod. Phys. 54, 235 (1982); C. Borgs and W. Janke, J. Phys. I France 2, 2011 (1992).
[7] J. Berges and I.-O. Stamatescu, Phys. Rev. Lett. 95, 202003 (2005).
[8] E.T. Tomboulis and A. Velytsky, Phys. Rev. D 72, 074509 (2005).
[9] T.R. Miller and M.C. Ogilvie, Nucl. Phys. B (Proc.
[10] A. Bazavov, B.A. Berg, and A. Velytsky, Int. J. Mod. Phys. A 20, 3459 (2005).
[11] B.A. Berg, *Markov Chain Monte Carlo Simulations and Their Statistical Analysis* (World Scientific, Singapore, 2004).
[12] N. Cabibbo and E. Marinari, Phys. Lett. B 119, 387 (1982).
[13] K. Fabricius and O. Hahn, *Phys. Lett.* B 143, 459 (1984); A.D. Kennedy and B.J. Pendleton, *Phys. Lett.* B 156, 393 (1985).
[14] B.A. Berg and T. Neuhaus, Phys. Rev. Lett. 68, 9 (1992).
[15] A.M. Ferrenberg and R.H. Swendsen, Phys. Rev. Lett. 61, 2635 (1988); 63, 1658 (1989).
[16] p.177 of [11].
[17] p.252 ff. of [11].
[18] A. Pelissetto and E. Vicari, Phys. Rep. 368, 549 (2002).
[19] F. Karsch and S. Stickan, Phys. Lett. B 488, 319 (2000); N. Alves, B.A. Berg and R. Villanova, Phys. Rev. B 43, 5846 (1991). Note that an extra factor 3/2 is used in the energy notation of the last article.
[20] p.484 of [4].
[21] C.M. Fortuin and P.W. Kasteleyn, *Physica* (Amsterdam) 57, 536 (1972); A. Coniglia and W. Klein, J. Phys. A 13, 2775 (1980).
[22] S. Fortunato, J. Phys. A 36, 4269 (2003).
[23] G. Boyd, J. Engels, F. Karsch, E. Laermann, C. Legeland, M. Lütgemeier, and B. Peterson, *Nucl. Phys.* B 469, 419 (1996).
[24] A. Hasenfratz and P. Hasenfratz, Phys. Lett. B 93, 165 (1980).
[25] J.D. Gunton and M. Droz, *Introduction to the Theory of Metastable and Unstable States* (Springer, Berlin, 1985).
[26] J.W. Cahn, Trans. Metall. Soc. AIME 242 (1968) 166.
[27] J.W. Cahn and J.E. Hilliard, J. Chem. Phys. 28, 258 (1958).
[28] A. Velytsky, Ph.D. thesis, Florida State University, Tallahassee, Florida, USA, 2004. UMI-31-37501. On the web at [http://etd.lib.fsu.edu](http://etd.lib.fsu.edu).
[29] O. Kaczmarek, F. Karsch, F. Zantow, and P. Petreczky, Phys. Rev. D 70, 074505 (2004); O. Kaczmarek and F. Zantow, Phys. Rev. D 71, 114510 (2005).
[30] P. Petreczky, Nucl. Phys. B (Proc. Suppl.) 140, 78 (2005); E. Laermann and O. Philipsen, Ann. Rev. Nucl. Part. Sci. 53, 163 (2003).
[31] J. Engels, F. Karsch, and T. Scheideker, *Nucl. Phys.* B 564, 303 (2000).
[32] Y. Deng, *Nucl. Phys. (Proc. Suppl.)* 9, 334 (1989); Engels, J. Fingberg, F. Karsch, D. Miller, and M. Weber, *Nucl. Phys.* B 252, 625 (1990).
[33] U.W. Heinz, AIP Conf. Proc. 739, 163 (2004); P.F. Kolb, Acta Phys. Hung. N.S. 21, 243 (2004); P.F. Kolb and U.W. Heinz, *Hydrodynamic Description of Ultrarelativistic Heavy-Ion Collisions in Quark Gluon Plasma 3*, R.C. Hwa and X.N. Wang, Editors (World Scientific, Singapore, 2004); nucl-th/0305084.
[34] R.D. Pisarski, *Phys. Rev.* D 62, 111501(R) (2000); A. Dumitru and R. Pisarski, *Phys. Lett.* B 504, 282 (2001).
[35] O. Scavenius, A. Dumitru and A.D. Jackson, *Phys. Rev. Lett.* 87, 182302 (2001).