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On the deconfining transition of SU(∞) gauge theory

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ABSTRACT

Monte Carlo simulations of Twisted Eguchi-Kawai models with asymmetric couplings have been used to investigate the large N deconfining phase transition of pure SU(N) gauge theories. Data with N=25,36,49 and 64 and variable asymmetry parameter ξ allow us to disentangle these theory's bulk first order transition from its physical deconfining transition. The SU(64) data is in good agreement with asymptotic freedom and gives a critical temperature of $T_c/\Lambda_E = 118 \pm 6$, or $T_c/\sqrt{\sigma} = 0.42 \pm 0.05$, which is close to the SU(3) ratio of deconfining temperature to the string tension. The data also suggest that the transition is first order.

I. INTRODUCTION.

In recent years there has been considerable effort in the study of the deconfinement transition in non-abelian gauge theories [1]. A detailed knowledge of the nature and properties of this transition is essential for a complete understanding of the confinement mechanism. Heavy ion collisions shall soon provide us with situations in which deconfinement can be studied in the laboratory - and it is important to know what the theory predicts. Furthermore, deconfinement transitions are relevant to the early history of the universe.

Lattice gauge theories provide a natural framework for a quantitative study of this phenomenon. It has been known for some time that the order of the transition depends on the gauge group. For the pure SU(2) gauge theory the transition is second order [2], while SU(3) shows strong first order behavior [3]: in conformity with theoretical expectations based on universality [4]. For SU(N) groups with $N \geq 4$, universality arguments do not predict the order unequivocally. There exist, however, several arguments pertaining to the order of the transition [5,6] : it is generally believed to be first order (see, however, Ref.[7]). Recently the SU(4) theory has been investigated : the results indicate a first order transition [8,9].

It is important to know the order and nature of the transition for higher values of N : particularly in the context of the large N limit. The large N approximation provides a valuable framework to study the dynamics of gauge fields. The study of deconfinement at $N=\infty$ would throw light on the mechanism of confinement ; and a comparison with the SU(3) behavior would provide a basis for examining the validity of the large N approximation itself.

Numerical investigations in large N theories have been made possible by the advent of Eguchi-Kawai models [10,11,13]. The crucial observation is that at $N=\infty$,field theories with an internal U(N) or O(N) symmetry become equivalent to matrix models living at a single site. Earlier, the Quenched Eguchi-Kawai model (QEK) [11] was studied at finite temperature [12] : evidence for deconfinement was found but the order of the transition was not clear. For numerical purposes, the Twisted Eguchi-Kawai (TEK) model [13] is more suitable. Both the TEK gauge theory at zero temperature and the TEK Chiral model in two dimensions [15] have been studied in detail. A finite temperature version of the TEK gauge theory was formulated in Ref.[16] using asymmetric twists. This hot TEK model has been studied in Ref.[17] and [18]. In Ref.[17] it was noted that it is rather difficult to study the physical aspects of deconfiment in this model. The reason is that for practical values of N_o, the time extent of the box, the first order bulk large N transition in the Wilson action interfered with the deconfinement transition. The former is not a deconfining transition, but involves a sharp change in the value of the string tension. For small N_o, this makes the confinement length larger than N_o, thus inducing a spurious deconfinement. This phenomenon persists upto $N_o = 4$ [18]. In principle, it is possible to separate the two transitions by using very large values of N_o ; thereby pushing the deconfinement transition to extremely weak couplings while the bulk transition remains at the same position. This, however , seems to be unpractical. There have been other studies on large N deconfinement using a slightly different hot TEK model [19] : we believe that their results are also plagued by the same problem.

A similar interference between bulk and deconfining transitions has been observed in SU(4) [8]. In this case, however, the bulk transition could be avoided in the standard way by adding a small adjoint piece to the Wilson action. This trick would not, however, work for large N, since a mixed action theory at $N=\infty$ is equivalent to a theory based on the standard Wilson action with a redefined coupling [20,21].

A different way of formulating a hot TEK model is to consider the standard symmetric twist with different couplings for the spatial and temporal plaquettes. This is equivalent to a field theory defined on an asymmetric lattice in a symmetric box. When the asymmetry parameter ξ (the ratio of the spatial lattice spacing to the temporal lattice spacing) is large, the physical temporal extent of the box is smaller than the spatial extent - thus simulating finite temperature effects. In a previous communication [20] we presented preliminary results of a Monte-Carlo study of the asymmetric lattice TEK model with N=16. It was found that for $\xi \geqslant 1.75$, the bulk transition disappears, but the Wilson line continues to show a discontinuous jump - indicating a first order deconfining transition freed forom the effects of the bulk transition. In this paper we present the detailed results of our work for N=16,25,49 and 64. The results for the larger values of N are qualitatively the same : though the critical value of ξ beyond which the action does not show any discontinuity seems to increase with N. Nevertheless we can always find a region of ξ for each N where the action is perfectly smooth and there is clear evidence for a first order deconfining transition. Our data for N=64 has strong evidence for scaling and the value of the physical deconfinement temperature is given by

$$T_c/\Lambda_E = 118 \pm 6$$

Using the string tension value of Ref. [14] one then has

$$T_c/\sqrt{\sigma} = 0.42 \pm 0.05$$

which is close to the value obtained in the SU(3) theory [3].

In Section II we recall the essentials of TEK models, some aspects of field theories on asymmetric lattices and the asymmetric coupling version of the hot TEK model. In Section III we present the results of our Monte Carlo simulations. Section IV contains analysis of the data and in Section V we make some concluding remarks.

II. THE MODEL

The Twisted Eguchi-Kawai model is defined by the partition function :

$$Z_{TEK} = \int \prod_{\mu} dU_{\mu} \exp\left(-S_{TEK}\right) \qquad (2.1)$$

where

$$S_{TEK} = -\beta \sum_{\mu>\nu} Z_{\mu\nu} \, Tr\left(U_{\mu} U_{\nu} U_{\mu}^{+} U_{\nu}^{+}\right) + h.c. \qquad (2.2)$$

U_{μ}'s are SU(N) matrices and $Z_{\mu\nu}$'s are constant elements of Z_N :

$$Z_{\mu\nu} = \exp\left(\frac{2\pi i}{N} n_{\mu\nu}\right) \qquad (2.3)$$

where $n_{\mu\nu}$'s are integers modulo N. Let L be an integer and $N=L^2$. Then with a symmetric twist, i.e.

$$n_{\mu\nu} = L \quad \text{for all} \quad \nu > \mu \qquad (2.4)$$

the above model is equivalent at $L=\infty$ to a SU(N) gauge theory defined in a symmetric box of size L [13]. The relationship between the link variables in the gauge theory and the reduced variables U_{μ} is summarised by the reduction prescription :

$$U_{\mu}(x) \rightarrow D(x)\, U_{\mu}\, D^{+}(x) \qquad (2.5)$$

where

$$D(x) = \prod_{\mu} \left(\Gamma_{\mu}\right)^{x_{\mu}}$$

and Γ_μ's are traceless SU(N) matrices forming a 't Hooft algebra :

$$\Gamma_\mu \Gamma_\nu = Z_{\mu\nu} \Gamma_\nu \Gamma_\mu \tag{2.6}$$

For any gauge invariant quantity $f(U_\mu(x))$ the following relation holds at $N=\infty$:

$$\left\langle f(U_\mu(x)) \right\rangle_{\substack{FIELD \\ THEORY}} = \left\langle f(D(x)\, U_\mu\, D^\dagger(x)) \right\rangle_{TEK} \tag{2.7}$$

where $\langle . \rangle_{TEK}$ denotes averaging in the ensemble defined by eqn. (2.1).

One way to introduce finite temperature in this model is to construct suitable asymmetric twists. Such twists, and the corresponding Γ_μ's were constructed in Ref.[16]. In this scheme, one can keep the temporal extent of the box in which the equivalent field theory is defined fixed, while the spatial extents go to infinity as N goes to infinity. Monte Carlo simulations of the hot TEK model with the twists of Klinkhamer and van Baal [16] have been performed in Refs.[17] and [18]. As discussed earlier, for practical values of N_o it is hard to decouple the bulk transition of this theory from the deconfinement transition. The possible values of N in such models are rather restrictive, rendering studies of higher N_o models unpractical.

An alternative way to simulate finite temperature in lattice gauge theories is to consider a theory defined in a symmetric box (i.e same number of lattice sites in all directions) but with different lattice spacings in the spatial and temporal directions. The asymmetry parameter is defined by :

$$\xi = a/a_\tau \tag{2.8}$$

where a and a_τ denote the lattice spacings in the spacelike and timelike directions respectively. In order to obtain the correct continuum limit one now needs different couplings for the spatial and temporal plaquettes. The action is given by :

$$S = \sum_x \left\{ \beta_\sigma \sum_{i \neq j = 1}^{3} P_{ij} + \beta_\tau \sum_{i=1}^{3} P_{0i} \right\} \tag{2.9}$$

where P_{ij} and P_{oi} denote the standard spacelike and timelike plaquette traces respectively. The continum limit is defined by :

$$a \to 0 \qquad\qquad \xi = \text{FIXED}$$

Let the above theory be defined in a box with L lattice sites in each direction. For sufficiently large ξ the physical time extent is much smaller than the spatial extent. In the limit :

$$a \to 0 , \quad L \to \infty$$
$$\xi , \; La_\tau = \text{FIXED}$$

the lattice model describes a finite temperature field theory with the physical temperature given by :

$$T = \xi/La \tag{2.10}$$

The absence of renormalisation of the velocity of light in the continum limit imposes relations between the two couplings $\beta_\sigma(a,\xi)$ and $\beta_\tau(a,\xi)$. In weak coupling one has :

$$\begin{aligned} \beta_\sigma(a,\xi) &= \frac{1}{\xi g_E^2(a)} + \frac{1}{\xi} c_\sigma(\xi) + O(g_E^2) \\ \beta_\tau(a,\xi) &= \frac{\xi}{g_E^2(a)} + \xi c_\tau(\xi) + O(g_E^2) \end{aligned} \tag{2.11}$$

where $g_E^{\,2}(a)$ is the "euclidean" bare coupling on a symmetric lattice. The functions $c_\sigma(\xi)$ and $c_\tau(\xi)$ have been calculated in weak coupling perturbation theory in Ref.[22].

A TEK model which is equivalent to the above asymmetric lattice theory at $N=\infty$ may be written down in a manner entirely analogous to the cold TEK model. The reduced action is :

$$S = -\beta_\sigma \sum_{i\neq j=1}^{3} Z_{ij}\, \mathrm{Tr}(U_i U_j U_i^\dagger U_j^\dagger) - \beta_\tau \sum_{i=1}^{3} Z_{0i}\, \mathrm{Tr}(U_0 U_i U_0^\dagger U_i^\dagger) + h.c. \tag{2.12}$$

where the $Z_{\mu\nu}$'s are the same as in eqns. (3) and (4). The correspondence between the variables in the field theory and in the reduced model is the same as in eqn.(5). In particular, the thermal Wilson line is given by :

$$\langle WL \rangle = \frac{1}{N}\langle \mathrm{Tr}\, U_0^L \rangle \tag{2.13}$$

This is the standard order parameter for the deconfinement transition. As usual at low temperatures a Z_N symmetry prevents this from acquiring a vacuum expectation value -

signifying confinement. At high temperatures this symmetry is broken and $\langle WL\rangle$ acquires a nonzero value. In our hot TEK model at extreme weak couplings U_0 is frozen near its vacuum value Γ_0 ; but since $\Gamma_0^L = 1$, $\langle WL\rangle \neq 0$. Note that L is the lowest integer for which Tr $U_0^L \neq 0$. This is just a special case of the fact that traces of all open lines vanish unless they run from one end of the box to the other. For standard reasons Tr $U_0 = 0$ in the strong coupling domain. Hence at some intermediate coupling $\beta_E = \beta_c$ there is a phase transition.

The expression for the 'total action', i.e. the average plaquette is :

$$\langle S\rangle = \frac{1}{N} \mathrm{Re} \left\langle \sum_{\mu>\nu} Z_{\mu\nu} \mathrm{Tr} \left(U_\mu U_\nu U_\mu^\dagger U_\nu^\dagger \right) \right\rangle \qquad (2.14)$$

$\langle S\rangle$ would show a discontinuous jump at a bulk first order phase transition.

III. MONTE CARLO RESULTS

We have performed Monte Carlo simulations of asymmetric coupling TEK model for N=16,25,36,49 and 64 and various values of ξ ranging from 1.5 to 4.0. The Metropolis scheme used for updating is discussed in Refs [13] and [17]. For each value of ξ, the functions $c_\sigma(\xi)$ and $c_\tau(\xi)$ were computed from the calculations performed in Ref.[22]. These functions are of the general form (for SU(N)) :

$$C_\sigma(\xi) = \frac{N^2-1}{6N} f_1(\xi) + 4N f_2(\xi) \qquad (3.1)$$

$$C_\tau(\xi) = \frac{N^2-1}{4N} g_1(\xi) + 4N g_2(\xi)$$

Curves of $c_\sigma(\xi)$ and $c_\tau(\xi)$ vs. ξ are given in Ref[22] for N=2 and N=3. From these curves the values of $f_1(\xi)$, $f_2(\xi)$, $g_1(\xi)$ and $g_2(\xi)$ were calculated for each ξ separately. Equations (11) were then used to compute $\beta_\sigma(a,\xi)$ and $\beta_\tau(a,\xi)$ in terms of $\beta_E = 1/g_E{}^2(a)$ and $c_\sigma(\xi)$ and $c_\tau(\xi)$. This ensures that in the weak coupling limit one is simulating finite temperature physics with the physical temperature given by eqn.(10). For each value of N and ξ, we then scanned over various values of β_E measuring the total action, the Wilson line and the energy density.

Let us first discuss the basic features of our results which are common for all values of N. The total action values show good agreement with the results of lowest order strong and weak coupling expansions at the respective ends.

The Wilson line is consistent with zero at strong couplings and makes a discontinuous jump at some intermediate coupling. The amount of the jump decreases slightly with increasing N. For small values of ξ the action, however, has a discontinuity at a place which coincides with the position where the Wilson line jumps for smaller N. This is the bulk first order phase transition. As discussed earlier, the bulk transition is driving a spurious deconfinement transition. The amount of the discontinuity of the action decreases as ξ increases. Beyond a certain value of $\xi = \xi_c$ (which increases with N) the action becomes continuous. The Wilson line continues to show a sharp jump - indicating a strong first order deconfinement transition not affected by any bulk transition.

The results for N=16 have been reported earlier in Ref.[20]. Here we summarise our results for the other values of N. In Fig.1 we show the total action and the Wilson line for N=25 in the range $0.3 < \beta_E/N < 0.4$ for ξ = 1.5,2.0,2.5 & 3.0. For all of these values of ξ the Wilson line jumps discontinuously from zero to about 0.4 at some value of the coupling. However, for ξ = 1.5 and 2.0 the action has a discontinuity at the same value of the coupling (This was checked by making hot and cold runs in the standarad fashion). Somewhere between ξ = 2.0 and 2.5 the discontinuity in the action disappears : the action is perfectly smooth for ξ = 2.5 and 3.0 while the Wilson line continues to jump. Fig.2 shows the same quantities for N =

36. At $\xi = 2.5$ the action is almost continuous; it is smooth at $\xi = 3.0$. However, for these values of ξ the coupling at which the Wilson line jumps is well into the strong coupling region. Fig.(3) shows the data for N = 49. Thre action becomes continuous only for $\xi \geqq 3.0$. For other values of ξ , the position of the discontinuity of the action coincides with the position where the Wilson line jumps. Fig.(4) shows the results of our N = 64 runs. Now there appears to be a slight discontinuity even at $\xi = 3.0$. However, at $\xi = 1.5$ the Wilson line jumps at $\beta_E/N = 0.3525$ which is clearly to the weak coupling side of the bulk transition which occurs at $\beta_E/N = 0.3425$. (Note that at this coupling a change of β_E/N by 0.0322 corresponds to a change of length scale by a factor of two, if we trust asymptotic scaling). This is the most reliable measurement of the deconfinement transition we could perform. For other values of ξ the bulk and deconfining transitions are not so clearly separated.

Each point in Figs.(1)-(4) represents block averages over typically 1000-2000 sweeps. Our typical resolution in β_E/N was 0.0025 as the figures indicate. The typical error in the total action was about 0.01, while that in the Wilson line was around 5%.

It may be noted that the value of ξ above which the action becomes smooth increases with N. We do not understand fully the mechanism by which a higher ξ makes the first order bulk transition disappear. One possibility is that as ξ is increased, the physical temporal extent of the box decreases. When this becomes smaller than the length scale of correlations causing the bulk transition, the latter is suppressed. Our data is roughly consistent with this possibility.

IV. DATA ANALYSIS

As our results indicate, it is extremely difficult to decouple the bulk transition from the physical deconfinement transition and at the same time push the latter into the weak coupling regime. In most cases the bulk transition occurs at the same place at which the Wilson line jumps. For high values of ξ the bulk transition smoothens out, but the Wilson line tends to jump in the strong coupling side. The only clear exception is the N=64, ξ=1.5 data. Here the bulk transition is still present, but deconfinement occurs in the weak coupling side. For other values of ξ the situation is not so clear. To determine whether this corresponds to deconfinement in the continuum theory, one must make sure that scaling has set it. Let T_c denote the physical deconfinement temperature. Then :

$$T_c = \frac{\xi}{La(\beta_c/N)} \tag{4.1}$$

where β_c is the critical euclidean coupling. If β_c is in the asymptotic scaling region, one would have :

$$\frac{T_c}{\Lambda_E} = \frac{\xi}{L}\left(\frac{11}{48\pi^2}\,\frac{N}{\beta_c}\right)^{51/121} \exp\left(\frac{24\pi^2}{11}\,\frac{\beta_c}{N}\right) \tag{4.2}$$

where Λ_E is the euclidean Λ paramter. Reversing the argument one could calculate T_c/Λ_E using (4.2) and see whether this is independent of ξ and L. Figure 5 shows T_c/Λ_E for various values of N as a function of ξ, using the data in Figs.1-4. While the others show a gross violation of scaling (as expected), the N=64 data does show some tendency towards scaling. This should be, however, viewed with caution, since for $\xi=2$ and 2.5 the critical coupling is in the region of the bulk transition, while for $\xi=3$ the bulk discontinuity has smoothened but β_c is at a much stronger coupling. However, scaling might have set in earlier than expected. As discussed the N=64, $\xi=1.5$ data is the most reliable measurement we could make. This gives our best estimate for T_c (assuming that scaling has set in) :

$$T_c/\Lambda_E = 118 \pm 6$$

To get some idea of the value of the deconfinement temperature in physical units let us use the string tension measurements performed on the symmetric lattice TEK model in Ref.[14] :

$$\sqrt{\sigma}/\Lambda_E = 280 \pm 20$$

This yields :

$$T_c/\sqrt{\sigma} = 0.42 \pm 0.05 \qquad (N=\infty)$$

This value of $T_c/\sqrt{\sigma}$ is very close to the SU(3) value [3] :

$$T_c/\sqrt{\sigma} = 0.50 \pm 0.05 \qquad (N=3)$$

To obtain more accurate numbers it is necessary to measure correlations of Wilson lines. The connected correlation is howvever down by $O(1/N^2)$ - this is the statement of factorisation- and impossible to measure in a EK model where factorisation is exact.

V. CONCLUDING REMARKS

The use of asymmetric coupling TEK model was crucial in making our investigations possible. In our model N=64 corresponds to a 8^4 lattice. To obtain a model using hot twists of Ref.[16] for a temporal extent N_o=8 one needs N=384 !

We have presented strong evidence in favor a first order deconfining transition in the pure SU(∞) gauge theory. The discontinuity in the Wilson line, however, decreases markedly with increasing N, i.e. an increasing box size. (Δ<WL> decreases from 0.4 at N=25 to 0.1 at N=64). This is very similar to what happens in the SU(3) theory, where Δ<WL> decreases from 0.33 at N_o=2 [3] to 0.08 at N_o=6 [23]. The latter is predominantly due to perimeter corrections to the Wilson line average which has to be divided out to extract the physical free energy of a quark. The free energy thus obtained is indeed independent of N_o, at least for large N_o. The situation at N=∞ is analogous. We have checked in our data that the decrease in the discontinuity of the Wilson line is roughly consistent with a perimeter effect. However, it has been discussed by Pisarski [7] that, in a certain sense ,a second order transition is "natural" in the N=∞ theory where one has a theory of free glueballs, and the transition is of the Hagedorn type. More detailed studies are required to clarify this point.

While we cannot conclude with certainity that scaling holds, there is definitely some indication of scaling in our N=64 data. If this is indeed true, this would indicate the existence of a continuum limit for TEK models -and for confinement at low temperatures. Clearly at lot more work has to be done to establish scaling in a definitive fashion and extract the physical deconfinement temperature accurately. This could be presumably done for higher values of N. Such studies are, however, extremely time-consuming with the Metropolis scheme we utilised. It is possible that the heat bath method for TEK models devised in Ref.[24] may make such investigations practicable.

If we accept that our N=64 data shows scaling it is significant that the deconfinement temperature in physical units is very close to the SU(3) value. This indicates that the confinement mechanism in N=3 and $N=\infty$ are very similar. In that case the large N approximation is a good approximation. It would be certainly worthwhile to continue the large-N program, particularly in the analytic front, where there is more chance of success compared to the N=3 theory.

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FIGURE CAPTIONS

Fig. 1: Total Action and Wilson Line for N = 25 and ξ = 1.5, 2.0, 2.5 & 3.0. The dots represent the Total Action; the crosses represent the Wilson Line.

Fig. 2: Total Action and Wilson Line for N = 36. The dots are data points for the Total Action, the crosses for the Wilson Line

Fig. 3: Total Action and Wilson LIne for N = 49. The dots are data points for the Total Action, the crosses for the Wilson Line.

Fig. 4: The Total Action and the Wilson Line for N = 64. The dots are data points for the Total Action, the crosses for the Wilson Line.

Fig. 5: The deconfinement temperature in units of the euclidean Λ parameter (calculated assuming that asymptotic scaling holds) versus the asymmetry parameter for the various values of N. (Scaling is valid if T_c/Λ_E does not depend on ξ).