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Authors
Greensite, J.
Halpern, M.B.

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Lawrence Berkeley Laboratory
UNIVERSITY OF CALIFORNIA

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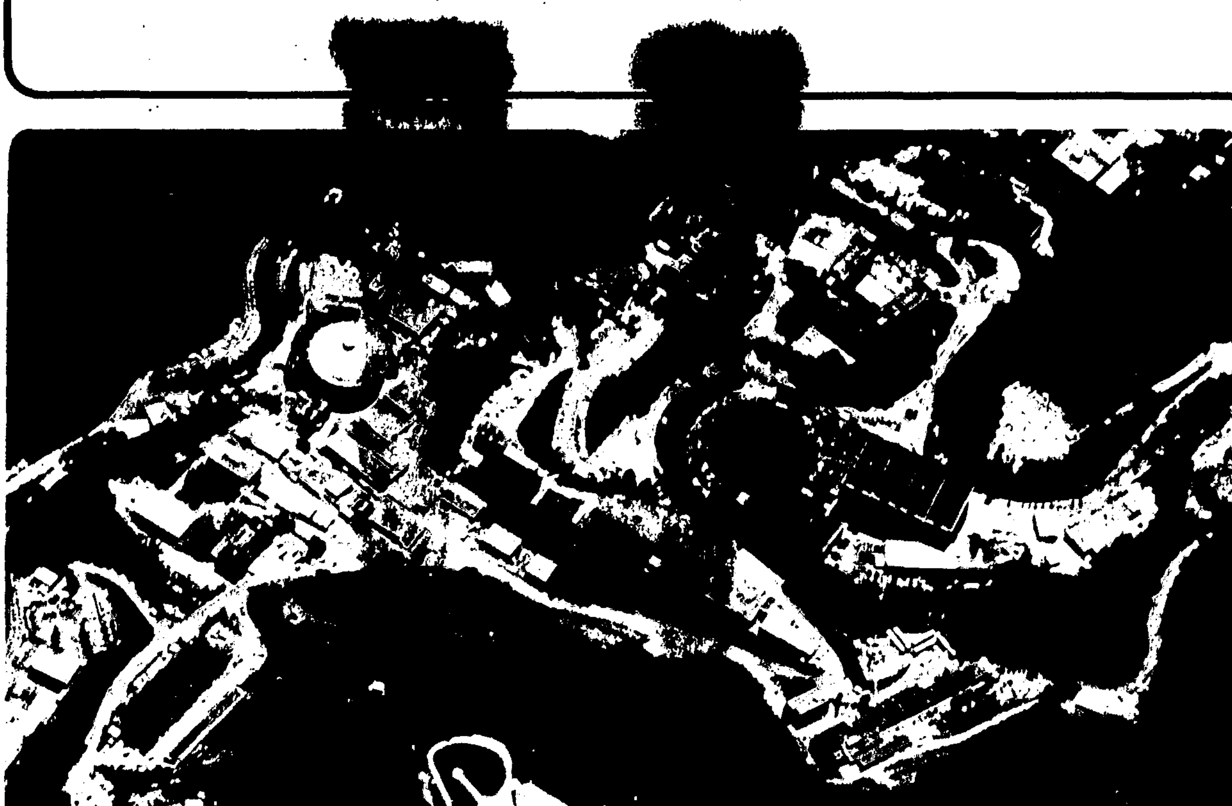

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SUPPRESSION OF COLOR SCREENING AT LARGE N*

J. Greensite

and

M. B. Halpern

Lawrence Berkeley Laboratory

and

Department of Physics

University of California

Berkeley, California 94720

ABSTRACT

In large N QCD, deconfinement by color screening is suppressed. The adjoint string tension is twice the fundamental string tension. Consequences for models of confinement are discussed, and a simple model of a confining large N master field is given.

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Consider an adjoint quark anti-quark pair in finite N QCD. If there is a (double) string between the quarks (confined state), it is expected that it is energetically favorable for the string to break (vacuum polarization), resulting in a state of free screened quark plus free screened antiquark.[1] In strong coupling Euclidean lattice gauge theory, the situation corresponds to Figs. (1a, b). A fixed time slice of the "sandwich" (Fig. 1a) is the bound state of the two quarks (connected by a double string). A fixed time slice of the "tube" (Fig. 1b) is the screened state (each quark free and wrapped in glue). The statement is that the sandwich contributes an area piece to the Wilson loop, while the tube contributes a perimeter piece. When the loop is large the sandwich is small and we see a perimeter law. This is color screening.

Our point in this report is that screening is suppressed in the large N limit. Computation of the diagrams in the figure gives

$$\langle Tr_A U[C] \rangle \sim N^2\{e^{-2\sigma_F A[C]} + N^{-2}e^{-4\sigma_F P[C]}\}, \qquad (1)$$

where Tr_A is the trace in the adjoint representation, and σ_F is the string tension for the fundamental quark. The first term is from the sandwich, the second from the tube. In the extreme large N limit, the tube is suppressed (color screening is suppressed), and we see an area law for the adjoint loop. In fact, comparing the terms in (1), we get an estimate for the screening length (at which the string will snap)

$$L \approx \left(\frac{\ln N^2}{2\sigma_F}\right)^{\frac{1}{2}}. \qquad (2)$$

A general argument for large N suppression of color screening deconfinement is available. Since (F = fundamental)

$$\mathrm{Tr}_A U[C] = \begin{cases} \mathrm{Tr}_F U[C]\mathrm{Tr}_F U^+[C], & (U(N)) \\ \mathrm{Tr}_F U[C]\mathrm{Tr}_F U^+[C] - 1, & (SU(N)) \end{cases} \tag{3}$$

it follows immediately from large N factorization that

$$< \mathrm{Tr}_A U[C] > \underset{N}{=} | < \mathrm{Tr}_F U[C] > |^2. \tag{4}$$

If $\sigma_{F(A)}$ is the string tension in the fundamental (adjoint) representation, Eq. (4) says that

$$\sigma_A \underset{N}{=} 2\sigma_F, \tag{5}$$

thus verifying our scenario above in generality. The relations (1) and (2) are also generic. Further, it is known that the topology of the leading contribution to the fundamental loop is planar, and order N. Equation (4) then states that the leading contribution to the adjoint loop has the topology of a sphere (with the adjoint loop at the equator), and order N^2. Note that the sandwich (Fig. 1a) is a sphere, while the tube (Fig. 1b) is a torus, and hence down in N.

In an extreme large N universe, we have seen that transitions: bound → screened are suppressed. By the same token, the number of free screened adjoint quarks is absolutely conserved (they cannot be produced or annihilated). In such a universe, created with no free screened adjoint quarks, none can arise, and it is fair to say all adjoint quarks are confined.

Our observations make it clear that conjectured confinement mechanisms having to do with the center of the group (Z_N fluxons,[2] spaghetti vacuum[3]) cannot survive at large N. Assume (fundamental) confinement is describable at all N in terms of a sum over fluxon configurations alone. Since the fluxons have no effect on the adjoint loop, the hypothesis is inconsistent with large N factorization, Eq. (4). A further set of ad hoc configurations X might be assumed to confine the adjoint quarks at large N, but X must be mysteriously correlated with the fluxons to produce (5). We find this unnatural. It is simpler to believe that the center of the group plays no role in large N confinement. The center of the group therefore stands to confinement roughly as instantons stand to the U(1) problems. It would be preferable to find a unified (all N) confinement mechanism. From a different direction, Lovelace has recently drawn similar conclusions about monopoles.[4]

Indeed the very idea of a group of configurations being necessary for large N confinement is presumably in contradiction with factorization. The standard lore is that, as N increases, some set Y of important configurations gradually shrinks to a single configuration, the master field.[5] We have argued here that the master field will confine adjoint quarks if it confines fundamental quarks. We now mention a simple toy master field that confines correctly and incorporates asymptotic freedom.

In Ref. [6], we have found an exact matrix equation (the quenched Langevin equation) for the QCD master field. In our approach, the master field is a translationally covariant function of 5N uniform random (quenched) momenta $\bar{p}^* = (p^*_{\mu a}, p^*_{5a})$ and a 4-vector N × N hermitean Gaussian random noise matrix η_μ^{*ab}. To zeroth order, the master field is

$$A_\mu^{*ab}(x) \approx e^{i(\bar{p}_a^* - \bar{p}_b^*)\cdot x}\,\eta_\mu^{*ab}$$

$$\cdot[i(p_{5a}^* - p_{5b}^*) + (P_a^* - p_b^*)^2]^{-1}$$

$$+ \text{ gauge terms} \qquad (6)$$

The prescription for the toy master field is to replace $(p_a^* - p_b^*)^2 \to \frac{11}{12\pi}(p_a^* - p_b^*)^2\, \ell n[1 + (p_a^* - p_b^*)^2/\mu^2]$ in Eq. (6). This gives rise to a two point function which corresponds to Richardson's potential.[7] The infra-red k^{-4} behaviour of the two point function is consistent with a truncated set of Schwinger-Dyson equations.[8]

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FIGURE CAPTIONS

Figure 1. Strong-coupling lattice diagrams responsible for (a) area law; and (b) perimeter law falloffs of the Wilson loop C in the adjoint representation.

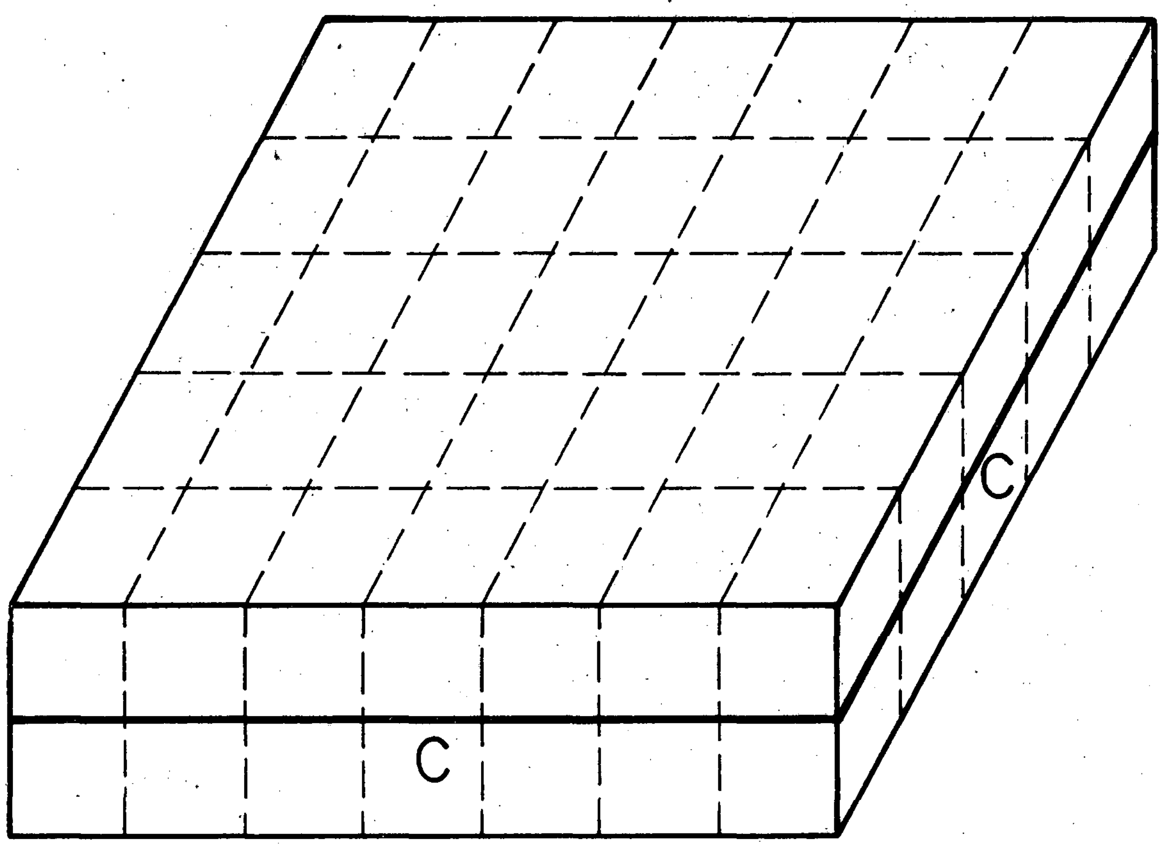

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FIGURE 1A

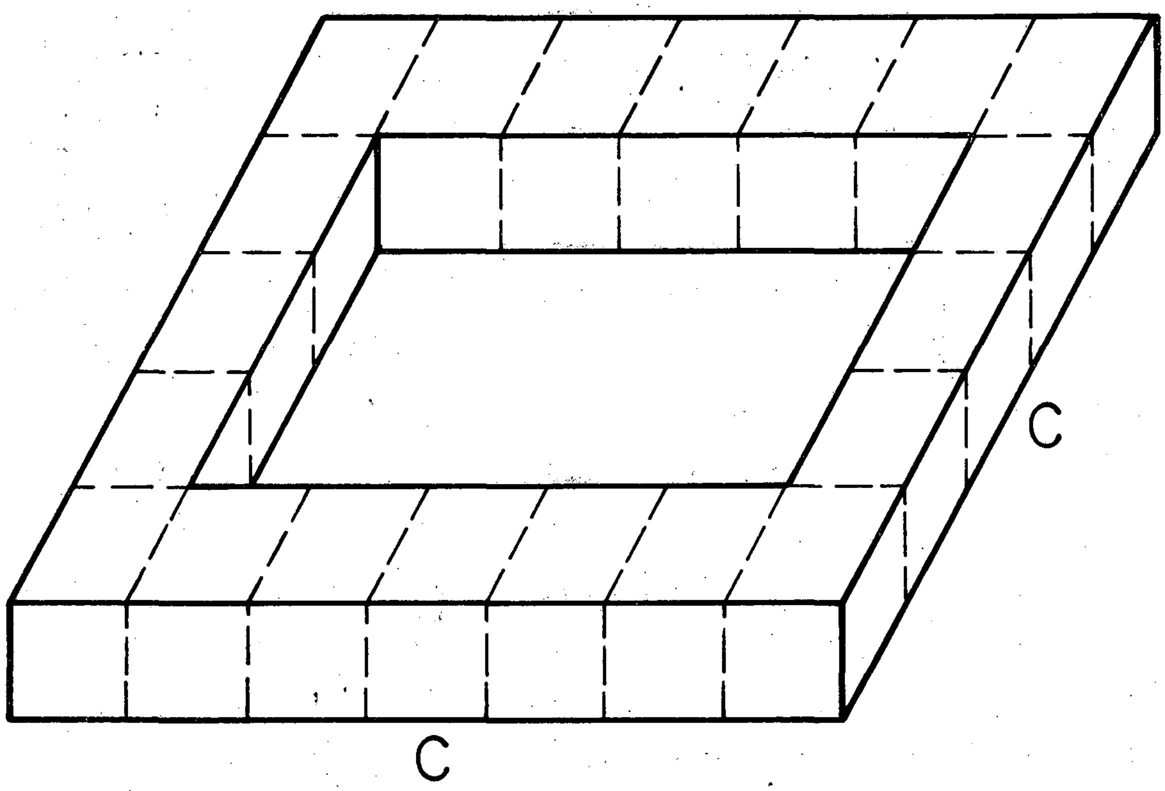

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FIGURE 1B

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LAWRENCE BERKELEY LABORATORY
UNIVERSITY OF CALIFORNIA
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