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INTERNATIONAL ATOMIC ENERGY AGENCY

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ABSTRACT

The Thermal Wilson Loop $< \mathrm{Tr}\ P \exp ig \int_0^\beta d\tau A_0(\tau, \vec{x}) >$, representing an order parameter for the gauge theory and expected to be zero in the confining phase, is perturbatively evaluated up to the $O(g^4)$ included for an SU(N) pure Yang Mills theory. This evaluation should be meaningful at high temperature, $\beta \to 0$. Its behaviour is discussed and a possible need for non perturbative instanton-like contributions is pointed out.

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In this paper we present an evaluation of the Thermal Wilson Loop based on perturbation theory. The Thermal Wilson Loop is defined to be

$$\langle \mathrm{Tr}\,\Omega(\vec{x}) \rangle = \langle \mathrm{Tr}\ P \exp ig \int_0^\beta d\tau\, \lambda^a A_0^a(\tau,\vec{x}) \rangle \tag{1}$$

where the functional average is taken over the gauge field $A_\mu^a(\tau, \vec{x})$ whose dynamics is here assumed to be described by a pure Yang-Mills Lagrangian without additional matter fields. We study the system in the continuum and we assume that it is placed in a heat bath at a physical temperature β^{-1}.* The integral over the Euclidean time τ in eq. (1) represents an integral over a closed path due to the periodicity conditions $A_\mu^a(0, \vec{x}) = A_\mu^a(\beta, \vec{x})$, P meaning the usual ordering. The Thermal Wilson Loop gives the exponential of minus the free energy of a static external particle[1,2,3], e.g. a heavy quark, transforming according to the representation λ^a, coupled to the gauge field at a temperature β^{-1}.

By defining the symbol Tr to be $\mathrm{Tr} = \frac{1}{d(R)}(\mathrm{trace})$ in such a way that $\mathrm{Tr}I = 1$, I being the identity and $d(R)$ the dimensionality of the given representation, we have for the limit of vanishing coupling constant $\lim_{g\to 0} < \mathrm{Tr}\,\Omega > = 1$. This is supposed to be also the high temperature limit $\beta \to 0$, since β is the only scale besides the renormalization scale μ of the quantum gauge theory and therefore the running coupling constant is $g = g(\beta\mu)$. It is speculated that a phase transition occurs[1,4,5] and that for β beyond a critical value, i.e. $\beta > \beta_c$, the Thermal Wilson Loop vanishes $< \mathrm{Tr}\,\Omega > = 0$, implying an infinite free energy for an isolated non singlet particle, which is consistent with confinement. Therefore the Thermal Wilson Loop $< \mathrm{Tr}\,\Omega >$ can be considered as an order parameter for this

* Sometimes in papers on lattice gauge theories the squared coupling constant g^2 is called "temperature". This is not so in our case.

phase transition. This picture has been confirmed by Monte-Carlo experiments in refs. [2], [6] and [7] and by semi-classical arguments in refs. [8] and [9]. From a more technical point of view a considerable amount of work has been done recently[10,11,12,13] in order to show, according to the Polyakov[14] suggestion, that all the logarithmic divergences coming from the perturbative expansion of a smooth contour Wilson Loop can be absorbed in the usual charge renormalization for the non abelian gauge field.

Here we have evaluated the first terms in the perturbative expansion in g of $< \mathrm{Tr}\,\Omega >$, which, according to what we have previously said, should be meaningful for studying the region around $\beta = 0$. We have used the finite temperature Feynman rules, working in the gauge $\partial_\mu A^a_\mu = 0$. Up to order g^4 included we have:

$$\begin{aligned}\mathrm{Tr}\,\Omega(0) = {} & 1 - g^2\,\mathrm{Tr}(\lambda_a\lambda_b)\, P\int\!\!\int_0^\beta d\tau_1 d\tau_2\, A^a_0(\tau_1,0)\,A^b_0(\tau_2,0) - \\ & - i\, g^3\,\mathrm{Tr}(\lambda_a\lambda_b\lambda_c)\, P\int\!\!\int\!\!\int_0^\beta d\tau_1 d\tau_2 d\tau_3\, A^a_0(\tau_1,0)\,A^b_0(\tau_2,0)\,A^c_0(\tau_3,0) + \\ & + g^4\,\mathrm{Tr}(\lambda_a\lambda_b\lambda_c\lambda_d)\, P\int\!\!\int\!\!\int\!\!\int_0^\beta d\tau_1 d\tau_2 d\tau_3 d\tau_4\, A^a_0(\tau_1,0)\,A^b_0(\tau_2,0)\,A^c_0(\tau_3,0)\,A^d_0(\tau_4,0)\ .\end{aligned} \tag{2}$$

The quantum average then gives:

$$\begin{aligned}\langle \mathrm{Tr}\,\Omega(0)\rangle = {} & 1 - g^2\,\frac{Q_0}{2}\int_0^\beta d\tau_1\int_0^\beta d\tau_2\, D'(\tau_1-\tau_2) + \\ & + \frac{1}{2}\left(g^2\,\frac{Q_0}{2}\int_0^\beta d\tau_1\int_0^\beta d\tau_2\, D(\tau_1-\tau_2)\right)^2 + g^4\,\mathrm{Tr}\left(\lambda_a[\lambda_b,\lambda_a]\lambda_b\right) \times \\ & \times P\int\!\!\int\!\!\int\!\!\int_0^\beta d\tau_1 d\tau_2 d\tau_3 d\tau_4\, D(\tau_1-\tau_3)\,D(\tau_2-\tau_4) = \\ = {} & 1 + g^2\,\mathcal{L}_2 + \frac{1}{2}\left(g^2\,\mathcal{L}^0_2\right)^2 + g^4\,\mathcal{L}_4\end{aligned} \tag{3}$$

where we have put $Q_0 = \sum_a \mathrm{Tr}(\lambda_a\,\lambda_a)$ i.e. $Q_0 = \frac{T(R)}{d(R)}\cdot r$, r being the dimension of the group, $D(\tau) = D_{00}(\tau,\vec{0})$ being the zero-zero component of the propagator in the Feynman gauge. $D'(\tau)$ and therefore $\mathcal{L}_2$ include the $O(g^2)$ radiative corrections. As it is seen, the third term of eq. (3) represents the exponentiation of the second term, which would be the whole contribution in an abelian theory. At the order g^4 it is proportional to $(\mathcal{L}^0_2)^2$, $\mathcal{L}^0_2$ being the lowest order one-gluon exchange contribution. The other fourth order piece coming from the third term in the expansion (2) combined with the three gluons interaction gives no contribution since it is proportional to

$$\int_0^\beta dx_0\int d\vec{x}\; P\int\!\!\int\!\!\int_0^\beta d\tau_1 d\tau_2 d\tau_3\left[(\partial_{\tau_3}-\partial_{\tau_2})+(\partial_{\tau_1}-\partial_{\tau_3})+(\partial_{\tau_2}-\partial_{\tau_1})\right]\times$$
$$\times\, D_{00}(\tau_1-x_0,\vec{x})\;D_{00}(\tau_2-x_0,\vec{x})\;D_{00}(\tau_3-x_0,\vec{x})$$

which is identically zero.

We have assumed a dimensional regularization scheme, where the space dimensions are continued from 3 to d. In this scheme it is expected that the terms we propose to evaluate are finite, since the usual logarithmic infinities appear at this order as factors of linear divergences which are removed by dimensional regularization. In particular $\mathcal{L}^0_2$ is zero since it is proportional to $\int_0^\beta d\tau D(\tau) = \frac{1}{(2\pi)^d}\int\frac{d^dk}{k^2} = 0$. The fourth term

PLEASE REPLACE p.9 by the following:

$$\langle \mathrm{Tr}\,\Omega \rangle = 1 + 2\pi^2 Q_0 \left\{ \sqrt{\frac{2C_2}{3}} \left(\frac{g^2}{8\pi^2}\right)^{3/2} + C_2 \left(\frac{g^2}{8\pi^2}\right)^2 \left[\ln \frac{g^2}{8\pi^2} + \ln \frac{2\pi^2 C_2}{3} + \frac{3}{2} \right] \right\} \quad . \qquad (10)$$

We remember that for $SU(N)$ $C_2 = N$ and Q_o depends on the representation according to which the static external particle transforms: in the fundamental representation $Q_o = \frac{1}{2}\frac{N^2 - 1}{N}$, in the adjoint one $Q_o = N$. It is seen that if the r.h.s. of eq. (10) is expressed as a function of $\lambda = g^2 N$, it has a finite limit for $N \to \infty$ at fixed λ.

One can easily check that the r.h.s. of eq. (10) is 1 plus a positive monotonic increasing function of g^2. The higher order terms are expected to have renormalizable ultraviolet divergences and as one of the consequences the coupling constant $g^2/8\pi^2$ will be promoted to a running coupling constant $\frac{g^2_R(\beta)}{8\pi^2} = (- \frac{11N}{3} \ln\beta\mu + \ldots)^{-1}$ in the r.h.s. of eq. (10) according to the general renormalization group arguments. Therefore we can read eq. (10) as giving the leading behaviour of $< \mathrm{Tr}\,\Omega >$ as a function of β for $\beta \to 0$.

Therefore we get from our evaluation a $< \mathrm{Tr}\,\Omega >$ which increases monotonically above 1 for increasing β, contrary to the expectation that $< \mathrm{Tr}\,\Omega >$ should decrease and become eventually zero for $\beta \to \beta c$. Of course, one can say that we have only computed few terms of the perturbative series.*) But we can try anyhow to make some general considerations, based on the leading term for $g \to 0$, represented by the contribution called $\mathcal{J}$ in eq. (7) to the term $\mathcal{L}_2$ of eqs. (3) and (5). It is clear that if instead of the dimensional regularization we introduce a cut-off Λ, for instance $\Lambda \sim a^{-1}$, a being a lattice spacing, we find for $\Lambda >> gM$:

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*) Terms of higher order in g could receive contributions from an infinite number of diagrams, due to infra-red divergences. See Ref.18 and references quoted there.

[18] A.D. Linde, Phys. Letters 96B, 289 (1980).

$\mathcal{L}_4$ represents the graph of Fig. 1. We have evaluated it by expressing

$$D(\tau) = \frac{1}{\beta(2\pi)^d} \sum_{n=-\infty}^{+\infty} e^{i\frac{2\pi}{\beta}n\tau} \int \frac{d^d k}{(\frac{2\pi}{\beta}n)^2 + k^2}$$

by performing first the integrals over τ and then these over k. As already noticed, the linear divergences disappear by dimensional regularization, and we end up with a sum over n proportional to

$$\beta^2 \left(\frac{2\pi}{\beta}\right)^{2d-4} \left(\pi^{d/2}\,\Gamma(1-d/2)\right)^2 \sum_{n=1}^{\infty} n^{2d-6} \quad .$$

Here we use the Riemann ζ_R function and

$$\lim_{d\to 3} \sum_{n=1}^{\infty} n^{2d-6} = \lim_{d\to 3} \zeta_R(6-2d) = -1/2 \quad .$$

The fourth term is evaluated this way to be

$$\mathcal{L}_4 = Q_0\, C_2\, \frac{1}{64\pi^2} \tag{4}$$

where $C_2 = N$ for $SU(N)$, i.e. the Casimir operator in the adjoint representation (we normalize the λ^a matrices in the fundamental representation so that trace $(\lambda_a \lambda_b) = \frac{1}{2}\,\delta ab$).

The term called $\mathcal{L}_2$ represents the contribution of a one gluon exchange including radiative corrections, as in Fig. 2 where the line with the blob is the complete propagator:

$$\mathcal{L}_2 = -\frac{Q_0}{2}\,\frac{\beta}{(2\pi)^d} \int \frac{d^d p}{Z_3\, p^2 + \Gamma^{(2)}(p)} \tag{5}$$

Here $\Gamma^{(2)}(p) \equiv \frac{-g^2}{(2\pi)^d}\,\frac{C_2}{\beta}\,W(p)$ represents the radiative corrections on the zero-zero component of the gluon propagator due to the gluon self-interactions and the interactions with the ghost field. $\Gamma^{(2)}$ is a function of $p = |\vec{p}\,|$ since in our case $p_0 = 0$. The essential new fact at finite temperature is that $\Gamma^{(2)}(p = 0) \neq 0$, representing the Debye screening effect[8,15]. We have therefore put $\Gamma^{(2)}(p = 0) = g^2 M^2$ and:

$$\Gamma^{(2)}(p) = g^2 M^2 - \frac{g^2}{(2\pi)^d}\,\frac{C_2}{\beta}\,\left(W(p) - W(0)\right)$$

where g^2M^2 is computed to be $M^2 = \frac{C_2}{3\beta^2}$. It is easy to check that only the zero-zero component of the inverse gluon propagator is different from zero at $p = 0$, and that the ghost inverse propagator including radiative corrections also remains zero at $p = 0$.

The quantity $W(p)$ is expressed as

$$W(p) = \sum_n \int_0^1 dx \int d^d k \left\{ \frac{(2d-1)k_0^2 + \vec{k}^2 + p^2\left(\frac{5}{2} - x(1-x)\right)}{\left[k_0^2 + \vec{k}^2 + p^2 x(1-x)\right]^2} - \frac{d}{k_0^2 + \vec{k}^2} \right\} \tag{6}$$

where $k_0 = \frac{2\pi}{\beta}\,n$. It is convenient to consider separately the contribution of the $n = 0$ mode, called $S(p)$, from the rest, $R(p)$ i.e. $W(p) - W(0) = S(p) + R(p)$. The reason is that, for $d = 3$, $S(p)$ is non analytic for $p^2 \to 0$, namely $S|_{d=3} = 2\pi^3 p$. We cannot simply expand the integrand of eq. (5) in powers of g^2, since the result of the integration is not analytic in the quantity g^2M^2 for $g^2M^2 \to 0$. We rather rewrite up to $O(g^4)$, with $Z_3 = 1 + g^2\delta Z_3$,

$$\int \frac{d^d p}{Z_3 p^2 + \Gamma^{(2)}(p)} = \int \frac{d^d p}{p^2 + g^2 M^2} + \frac{g^2}{(2\pi)^d} \frac{C_2}{\beta} \int \frac{d^d p \, S(p)}{[p^2 + g^2 M^2]^2} + \tag{7}$$

$$+ \frac{g^2}{(2\pi)^d} \frac{C_2}{\beta} \int \frac{d^d p}{p^4} \left(R(p) - p^2 \delta Z_3 \right) = \mathcal{J} + \mathcal{S} + \mathcal{R} \, .$$

The first term gives, for $d = 3$, $g^2 \mathcal{J} = -2\pi^2 g^2 [\Gamma^{(2)}(p=0)]^{1/2}$, giving rise to a $O(g^3)$ contribution to the Thermal Wilson Loop. Actually, since we want to compute the $O(g^4)$ term, we have to ask ourselves if the mass term $\Gamma^{(2)}(p = 0)$ receives also contributions of the order $O(g^3)$. This happens to be the case, since in computing the zero-zero component of the gluon bubble of Fig. 3 we can take into account the mass term $g^2 M^2$ at the lowest order in one of the two branches (it is seen that if one of the two branches is the zero-zero component the other one is a space-space component of the propagator. Moreover the ghost term and the term due to the four gluons interaction do not give corrections of this kind.) This means adding to $W(0)$ an extra term

$$\delta W(0) = \sum_m \int_0^1 dx \int d^d k \; k^2 \left\{ \frac{1}{[k^2 + x g^2 M^2]^2} - \frac{1}{(k^2)^2} \right\} \, .$$

Again, the $n = 0$ mode is non analytic in $g^2 M^2$ and gives for $d = 3$ the leading correction $\delta W(0) = -2\pi^2 gM$ with M at the lowest order. This gives $\Gamma^{(2)}(p=0) = -\frac{g^2}{(2\pi)^d} \frac{C_2}{\beta} (W(0) + \delta W(0))$ and finally:

$$\mathcal{J} = - \left(2\pi^2 g \sqrt{\frac{C_2}{3}} + g^2 \frac{\pi}{4} C_2 \right) \frac{1}{\beta} \, . \tag{8}$$

The second term at the r.h.s. of (7), called $\mathcal{S}$, is expressed as an integral which is infra-red logarithmically divergent at $d = 3$ for $g^2 M^2 \to 0$. We expect therefore $\mathcal{S} \sim -g^2 \, \ln g^2 M^2$. It is also ultraviolet logarithmically divergent at $d = 3$, but this divergence cancels, as expected, with the similar one of the last term $\mathcal{R}$ (the piece proportional to δZ_3 being zero by dimensional regularization).

We find:

$$\mathcal{S} = s(d) \, (g^2 M^2)^{d-3} \, \Gamma(3-d) \frac{g^2}{(2\pi)^d} \frac{C_2}{\beta}$$

$$\mathcal{R} = \left[r(d) \, \beta^{6-2d} \, 2 \zeta_R(6-2d) \, \Gamma(3-d) + \frac{5}{2} \pi^4 \right] \frac{g^2}{(2\pi)^d} \frac{C_2}{\beta}$$

where $s(d)$ and $r(d)$ are regular as $d \to 3$:

$$s(d) = 4\pi^4 \left(1 - (3-d) \ln \pi \right) + O((3-d)^2)$$

$$r(d) = 4\pi^4 + (3-d) \pi^4 \left(\frac{9}{2} - 12 \ln \pi - 16 \ln 2 \right) + O((3-d)^2) \, .$$

Since $2 \zeta_R(6 - 2d) = -1 - (6 - 2d) \ln 2\pi + O((3 - d)^2)$ we find finally for $d = 3$:

$$\mathcal{S} + \mathcal{R} = - \frac{g^2}{(2\pi)^3} \frac{C_2}{\beta} \pi^4 \left[2 + 4 \ln\left(\frac{C_2 2\pi^2}{3} \right) + 4 \ln \frac{g^2}{8\pi^2} \right] \, . \tag{9}$$

We can then sum up the various contributions to the r.h.s. of eq. (3), where $\mathcal{L}_2^0 = 0$, $\mathcal{L}_4$ is given in eq. (4) and $\mathcal{L}_2$ is obtained from eqs. (5), (7), (8) and (9), and so we end up, expressing everything in terms of $g^2/8\pi^2$, with:

$$\langle \mathrm{Tr}\,\Omega \rangle = 1 + 2\pi^2 Q_0 \left\{ \sqrt{\frac{2C_2}{3}} \left(\frac{g^2}{8\pi^2}\right)^{3/2} + C_2 \left(\frac{g^2}{8\pi^2}\right)^2 \left[\ln \frac{g^2}{8\pi^2} + \ln \frac{2\pi^2 C_2}{3} + \frac{3}{2} \right] \right\} \quad . \qquad (10)$$

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One can easily check that the r.h.s. of eq. (10) is 1 plus a positive monotonic increasing function of g^2. The higher order terms are expected to have renormalizable ultraviolet divergences and as one of the consequences the coupling constant $g^2/8\pi^2$ will be promoted to a running coupling constant $\frac{g^2_R(\beta)}{8\pi^2} = (-\frac{11N}{3} \ln\beta\mu + \dots)^{-1}$ in the r.h.s. of eq. (10) according to the general renormalization group arguments. Therefore we can read eq. (10) as giving the leading behaviour of $<\mathrm{Tr}\,\Omega>$ as a function of β for $\beta \to 0$.

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$$g^2 \mathcal{L}_2 \sim -\left(g^2 \beta\Lambda - \frac{\pi}{2} g^3 \sqrt{\frac{C_2}{3}} \right) \frac{Q_0}{(2\pi)^2} \quad .$$

In this computation the term $-g^2\Lambda\beta$, being negative and increasing in absolute value for β increasing, appears to be in good shape for tuning the expected behaviour of $<\mathrm{Tr}\,\Omega>$. This term however is not there by using the dimensional regularization, and moreover it does not fit with the requirement of the renormalization group, according to which powers of β can only appear in terms containing a factor $\exp(-\mathrm{const.}/g^2)$, i.e. terms which cannot be seen in a perturbative expansion around $g^2 = 0$.

We observe that terms of this kind are indeed given by the finite temperature instanton contribution to $<\mathrm{Tr}\,\Omega>$, which has the form

$$\int \frac{d\rho}{\rho^5} \left(\frac{8\pi^2}{g^2(\rho)}\right)^N e^{-8\pi^2/g^2(\rho)} \, e^{-C\rho^2/\beta^2} f(\rho/\beta)\, \beta \int d\vec{x} \left(\mathrm{Tr}\,\Omega_{c\ell}(\vec{x},\rho,\beta) - 1 \right)$$

where $f(\rho/\beta)$ is a function behaving like a power and $\Omega_{c\ell}(\vec{x}, \rho, \beta)$ is the classical instanton configuration for $P \exp i \int_0^\beta A_o^a \lambda^a(\vec{x}, \tau) d\tau$. Let us notice the gaussian factor $\exp - C\rho^2/\beta^2$, C being a constant, which renders meaningful the instanton computation at high temperature [8]. This instanton contribution has been numerically evaluated in ref. [9] and it appears to play the role of the term previously discussed, namely it is negative and increasing in absolute value as a power of β. It could therefore be an essential ingredient in order to trigger the expected phase transition to the confining regime [8], [16], [17].

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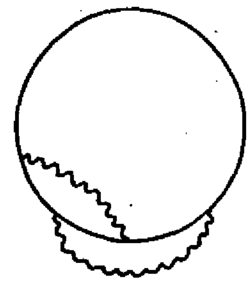

fig. 1

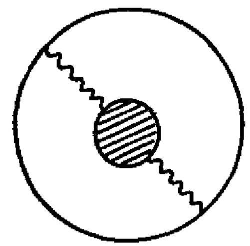

fig. 2

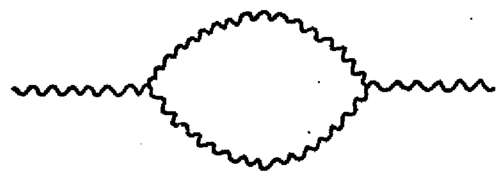

fig. 3

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