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FINITE TEMPERATURE SPINOR PREGOMETRY

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FINITE TEMPERATURE SPINOR PREGOMETRY *

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Abstract

We obtain the Einstein and Yang-Mills actions as induced actions. Free fermion fields are considered as the only fundamental pregeometric objects. We study in detail the temperature behaviour of the induced Newton and gauge coupling constants. Both of them are seen to decrease as the temperature grows.

Riassunto

Si ottengono le azioni di Einstein e Yang-Mills come azioni indotte. Campi fermionici liberi sono considerati come i soli oggetti pregeometrici fondamentali. Si studia in dettaglio il comportamento con la temperatura delle costanti di Newton e di gauge indotte. Si vede che entrambe diminuiscono al crescere della temperatura.

It is generally believed that spinor fields are the elementary constituents of matter. The forces among them are mediated by the gauge fields. All of these fields may be taken as fundamental in the sense that they can be given a dynamical role in the lagrangian. Alternatively in view of a theory in which matter and forces are expressed in terms of a single type of fundamental fields one can try to interpret the gauge fields as bound states of spinor fields⁽¹⁾. In this approach the action for the spinor fields endowed with a local symmetry gains the kinetic term of the relevant gauge field at the quantum level. The gauge fields result induced by the quantum fluctuations of the spinor fields. This latter attitude is quite similar to that adopted by Schwinger⁽²⁾ in order to obtain classical electrodynamics from the Dirac lagrangian minimally coupled to an external non-dynamical electromagnetic field and by Sakharov⁽³⁾ in order to obtain classical general relativity from a quantum matter generally covariant action.

Sakharov's proposal for induced quantum gravity has been developed by many authors⁽⁴⁾ but most of them use both fermions and gauge fields as fundamental objects.

We adopt a more "democratic" point of view in which gravity and Yang-Mills fields are treated on the same footing in the sense that both of them are induced by quantum fluctuations of matter fields⁽⁵⁾.

The central problem of this approach is the evaluation of the induced coupling constants.

As far as gravity is concerned, it is particularly interesting to study the behaviour of the induced Newton constant in extreme conditions as they can be realized e.g. in the early universe.

This motivates the finite temperature treatment of the whole

problem (6). We shall therefore study the dependence of the induced cosmological, Newton's and gauge coupling constants at finite temperature in a curved background.

We start from the gauge and generally covariant spinor lagrangian *

$$\mathcal{L} = \sum_f \bar{\Psi}_f (i \not{e}^{\mu\nu} \gamma_\mu \nabla_\nu - m) \Psi_f + h.c. \quad (1)$$

where Ψ_f is the fundamental representation of the local $SU(N_c)$ colour group, N_f is the number of flavours **, $e^{\mu\nu}$ is the vierbein field and $\nabla_\mu \equiv \partial_\mu - A_\mu(x) + \omega_\mu(x)$ being $A_\mu(x)$ the $N_c^2 - 1$ gauge potentials of $SU(N_c)$; $\omega_\mu(x)$ are the spin connections.

After functional integration over the fundamental fields the effective action depends only on the gauge fields and reads

$$W[A_\mu(x), e^{\mu\nu}(x)] = -\frac{1}{2} \text{Tr} \ln \Delta \quad (2)$$

where $\Delta \equiv -\square_x + m^2 + \frac{R}{4} + F_{\mu\nu} Z^{\mu\nu}$

and $\square_x$ is the generally covariant d'Alembertian on spinors, $F_{\mu\nu} = \partial_\mu A_\nu - \partial_\nu A_\mu + [A_\mu, A_\nu]$,

$Z^{\mu\nu} \equiv \frac{1}{4} [\gamma^\mu, \gamma^\nu]$; R is the Ricci scalar of the static background manifold. We assume the latter to have the topological structure $S^1 \times M^3 \equiv M$ with $\partial M = \emptyset$ so that we shall deal with thermal Green functions and avoid boundary terms. As usual we work on the euclidean section of the manifold so that the metric can be given the general form:

$$ds^2 = g_{\mu\nu}^E(x) dx^\mu dx^\nu = dt^2 + g_{ij}(\vec{x}) dx^i dx^j$$

* We use natural units $\hbar = c = k_B = 1$

** For the sake of simplicity we shall assume the same mass for all the different flavours.

In order to calculate the effective action we shall use the Schwinger - De Witt proper time Δ formalism (7) so that:

$$W = \frac{i}{2} \int_0^\infty \frac{ds}{s} \mathcal{H}(x, x; s) \quad (3)$$

where $\mathcal{H}$ is the heat kernel evaluated in the coincidence limit and here contains terms at zero and at finite temperature as a consequence of the global structure of this specific manifold. The explicit expression of $\mathcal{H}$ can be represented as an asymptotic expansion in powers of S for S going to zero. The series contains finite and infinite terms. The contribution due to the finite terms is negligible at low energies with respect to the mass scale while the divergent terms, once regularized, supply the induced quantities.

The finite temperature heat kernel can be expressed as a sum of zero temperature images antiperiodic in the imaginary time (8) so that:

$$W = W^{(0)} + W^{(R)} = -\frac{1}{2(4\pi)^2} \sum_k B_k \int_0^\infty ds s^{\frac{k-3}{2}} e^{-m^2 s} \left[1 + 2 \sum_{l=1}^\infty (-1)^l e^{-\beta^2 m^2/4s} \right] \quad (4)$$

where β , the finite length of the time axis as expressed by S^1 , is interpreted as an inverse temperature $1/T$ and $B_k \equiv \int_M d^4x \sqrt{g} b_k$ are the coefficients of the heat kernel expansion. The divergent part is formed of only the B_0, B_2, B_4 terms which read (9):

$$B_0 = 4 N_f N_c \int_M d^4x \sqrt{g}$$

$$B_2 = -\frac{1}{2} N_c N_f \int d^4x \sqrt{g} R$$

$$B_4 = \int d^4x \sqrt{g} \left\{ \frac{N_f}{6} F_{\mu\nu} F^{\mu\nu} + \frac{1}{2} \left[\frac{1}{2} (D_\mu F_\nu - F_{\mu\nu})^2 - \frac{1}{12} F_{\mu\nu}^3 \right. \right. \\ \left. \left. + \frac{1}{180} (R_{\mu\nu\alpha\beta} R^{\mu\nu\alpha\beta} - R_{\mu\nu} R^{\mu\nu} + 6 R^{\mu\nu}{}_{;\mu\nu}) - \frac{1}{6} (D_\mu R + \Sigma^{\mu\nu} F_{\mu\nu})^2 \right] \right\},$$

where we indicate with ∇ the unit spinor in the space of the internal quantum numbers of the fields involved. For the purposes of this paper it is quite reasonable to retain only the first term appearing in the expression for B_4 . The numerical factors containing N_f, N_c are the result of the trace on the internal quantum numbers. It is remarkable that the gauge field appears in the form of a Yang-Mills action only in B_4 which is the term containing all the quadratic combinations in the curvature. This is consistent with the geometrical meaning of the gauge field strength which is the curvature of the principal fibre bundle having $A(x)$ as a connection⁽¹⁰⁾ and with the quadratic character of the Yang-Mills action.

In (4) we can recognize the zero temperature contribution in the first term while the series summed over m is the finite temperature part of the action. The zero temperature induced constants are, after regularization with an ultraviolet cut off Λ on the divergent S -integration:

$$\frac{2\lambda}{16\pi G_N} = \frac{2N_c N_f}{(4\pi)^2} \left[\frac{\Lambda^4}{2} - m^2 \Lambda^2 + \frac{m^4}{2} \ln \frac{\Lambda^2}{m^2} - \frac{m^4}{2} (\gamma - 3/2) \right] \quad (5a)$$

where Λ is the cosmological constant

$$\frac{1}{16\pi G_N} = \frac{N_c N_f}{6(4\pi)^2} \left[\Lambda^2 - m^2 \ln \frac{\Lambda^2}{m^2} + (\gamma - 1) m^2 \right] \quad (5b)$$

$$\frac{1}{g^2} = \frac{N_f}{12(4\pi)^2} \left[\ln \frac{\Lambda^2}{m^2} - \gamma \right] \quad (5c)$$

Neglecting non-leading terms we can write (5) as

$$\frac{2\lambda}{16\pi G_N} \simeq \frac{N_f N_c}{(4\pi)^2} \Lambda^4 \quad (6a)$$

$$\frac{1}{16\pi G_N} \simeq \frac{N_f N_c}{6(4\pi)^2} \Lambda^2 \quad (6b)$$

$$\frac{1}{g^2} \simeq \frac{N_f}{12(4\pi)^2} \ln \frac{\Lambda^2}{m^2} \quad (6c)$$

From (6b) and (6c) we can easily recover a relation between G_N and g

$$\frac{1}{g^2} \simeq \frac{N_f}{12(4\pi)^2} \ln \left(\frac{6\pi}{G_N m^2 N_f N_c} \right) \quad (7)$$

(7) is analogous to the $\alpha_s G_N$ relation⁽¹¹⁾ obtained when pregeometry is combined with QED.

Relationships between different coupling constants are a typical result in pregeometric theories. It is remarkable that the sign of G_N in this case is uniquely defined and is positive as it must be. We recall that in scalar pregeometry the sign of G_N depends on the value of the non minimal coupling constant⁽⁶⁾. Moreover positivity of G_N rests on not having used as fundamental the gauge fields which give it a negative

contribution (12).

The finite temperature part of (4) reads

$$W^{(B)} = \frac{2}{(4\pi)^2} \sum_k B_k \sum_{l=0}^{\infty} (-1)^l \left(\frac{\beta^2 m^2}{4\pi^2} \right)^{\frac{k}{2}-1} \left| K_{\frac{k}{2}-2}(\beta m m) \right|, \quad (8)$$

where $K_\nu(x)$ are the modified Bessel functions defined according to

$$\int_0^\infty dx x^{\nu-1} e^{-\frac{x}{2} - \gamma x} = 2 \left(\frac{\alpha}{\gamma} \right)^{\nu/2} K_\nu(2\sqrt{\alpha\gamma})$$

and having made the substitutions

$$\gamma = \frac{k}{2} - 2$$

$$\gamma = m^2$$

$$\alpha = \frac{\beta^2 m^2}{4}$$

In (8) terms with $k > 4$ are negligible at curvatures low with respect to m^2 . Physically meaningful deviations from the zero temperature values of the induced coupling constants are expected to appear only at very high temperature. This justifies investigating the small β behaviour⁽¹³⁾ of (8):

$$W^{(B)} = \frac{1}{8\pi^2} \left\{ B_0 \left[-\frac{28\pi^4}{360\beta^4} + \frac{\pi^2 m^2}{6\beta^2} + \frac{m^4}{4} \left(\beta m \frac{\beta m}{4\pi} + \gamma - \frac{3}{4} \right) \right] + \right. \\ \left. + B_2 \left[\frac{\pi^2}{6\beta^2} + \frac{m^2}{2} \left(\beta m \frac{\beta m}{4\pi} + \gamma - \frac{1}{2} \right) \right] \right. \\ \left. - \frac{B_4}{2} \left(\beta m \frac{\beta m}{4\pi} + \gamma \right) \right\}. \quad (9)$$

From (9) we derive the temperature dependence of the induced coupling constants:

$$\frac{2\lambda(\beta)}{16\pi G(\beta)} = \frac{2\lambda}{16\pi G_0} + \frac{N_F N_c}{2\pi^2} \left[-\frac{7\pi^4}{90\beta^4} + \frac{\pi^2 m^2}{6\beta^2} + \frac{m^4}{4} \left(\beta m \frac{\beta m}{4\pi} + \gamma - \frac{3}{4} \right) \right] \quad (10a)$$

$$\frac{1}{16\pi G(\beta)} = \frac{1}{16\pi G_0} + \frac{N_F N_c}{24\pi^2} \left[\frac{\pi^2}{6\beta^2} + \frac{m^2}{4} \left(\beta m \frac{\beta m}{16\pi^2} + 2\gamma - 1 \right) \right] \quad (10b)$$

$$\frac{1}{g^2(\beta)} = \frac{1}{g^2} - \frac{N_F}{24\pi^2} \left(\beta m \frac{\beta m}{4\pi} + \gamma \right) \quad (10c)$$

The interesting features of (10b) and (10c) are exhibited taking into account only the leading terms. It results that

$$g^2(\tau) \simeq \frac{g^2}{1 + \frac{g^2 N_F}{24\pi^2} \beta m \frac{I}{m}} \quad (11a)$$

$$G(\tau) \approx \frac{G_\nu}{1 + \frac{\pi M_\rho^2 N_c}{g} G_\nu \tau^2} \quad (11b)$$

which decrease at very high temperature.

As far as $g(\tau)$ is concerned this behaviour should have been expected on the ground of the asymptotically free nature of Yang-Mills theory. It is less trivial recognizing ^{the} analogous behaviour for $G(\tau)$ although this is consistent with finite temperature scalar pregeometry ⁽⁶⁾. The decreasing of G with increasing T has been guessed by Zee in the framework of a particular toy-model ⁽¹⁴⁾. Comparison of (11a) and (11b) shows that the asymptotically free regime is reached at different temperatures. For the gauge coupling constant this occurs when T becomes large with respect to m , the mass of the fundamental fermion field, while G does not contain the memory of this mass and becomes small only when T becomes large with respect to the Planck mass $M_P \equiv G_\nu^{-1/2}$.

Therefore, $m \ll M_P$ even in grand unified theories, we may conclude that the asymptotic freedom regime is reached by $g(\tau)$ at temperatures much lower than those required by $G(\tau)$.

In different models $G(\tau)$ is induced by a symmetry breaking ⁽⁴⁾ of some matter field and $G(\tau) \propto \langle \phi_0(\tau) \rangle^2$ being $\phi_0(\tau)$ the temperature dependent order parameter ⁽¹⁵⁾. The value and also the sign of $G(\tau)$ can be strongly modified by the occurrence of phase transitions leading to the possibility of vanishing $\langle \phi_0(\tau) \rangle$. In the present work $G(\tau)$ is induced by the quantum fluctuations of a free matter field and no phase transitions can take place. Therefore the positive sign of G is preserved and no significant deviation from its zero temperature value occurs unless the temperature reaches very high values as can be inferred from (11b).

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