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INDUCED QUANTUM TORSION

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INDUCED QUANTUM TORSION *

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ABSTRACT

We study pregeometry in the framework of a Poincaré gauge field theory. The Riemann-Cartan space-time is shown to be an "effective geometry" for this model in the low energy limit. By using Heat Kernel techniques we find the induced action for curvature and torsion. We obtain in this way the usual Einstein-Hilbert action plus an axial Maxwell term describing the propagation of a massless, axial vector torsion field.

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General Relativity provides us with an excellent description for the large scale gravitational phenomena. Unfortunately Einstein's theory results to be unable to give a satisfactory account of the gravitational interaction at the sub-nuclear level. The still unresolved incompatibility between General Relativity and Quantum Mechanics could be skipped by assuming that General Relativity is only the low energy, or long wave-length limit of some more foundamental theory. In particular Sakharov showed in 1968 how the Einstein-Hilbert action for the spacetime geometry can be derived from the action of a general quantum matter field [1]. This approach is commonly referred to as Induced Gravity or Pregeometry; sometime the Sakharov's original idea is also called metrical elasticity because of some analogy with the theory of elasticity in solids [2]. The main problem is, in this framework, the evaluation of the induced Newton constant G_N in terms of the quantities characterizing the pregeometric matter action . In this regard, it seems that only massless spinor fields endowed with some internal gauge symmetry could provide a viable mechanism to generate G_N [3),4)]. On the other hand, we know from the Poincare' gauge theory of gravitation that the spin current density, associated to the spinor matter acts as the source of the torsion field just like the energy-momentum tensor is the source of the curvature [5]. In this way the two foundamental features of matter, i.e. spin and mass-energy are both coupled to the spacetime geometry which results to be of the Riemann-Cartan type.

In this letter we will show that whenever spinors are assumed as foundamental pregeometric fields, the corresponding induced action contains torsion terms besides the usual Einstein-Hilbert part. This result is obtained by choosing as an affine connection precisely the connection defined by the spinor structure itself through the metricity condition [6]. We shall discuss this point in the first part of this note. In the second part the Heat-Kernel expansion method [7] will be used to recover the induced action [8]. Spectral geometry in the Riemann-Cartan space-time has been investigated by several authors [9),10),11),12)]. We think that the main technical problem in this framework, that is the very large number of functionally independent invariants of the curvature and torsion entering the "Hamidew" coefficients has been definitly clarified in ref.12). The key point is that the Dirac field in U^4 interacts only with the axial vector part of the torsion tensor. This follows from the fact that the spin density current coupled to the contortion part of the geometry is totally skew-symmetric, therefore it selects among the irreducible components of the torsion tensor the axial vector piece only [13),14)]. The final result is that one can formally treat the Dirac field in U^4 like it were in a purely Riemannian geometry plus an additional interaction with an axial gauge field. Such an extra U(1) chiral gauge symmetry strongly reduces the number of admissible terms in the "Hamidew" coefficients of the Heat-Kernel expansion.

Physical quantities are expressed in natural units $\hbar = c = 1$. We adopt the following index notation ; $\lambda, \mu, \nu, \dots$ = coordinate indeces; a,b,c,...= Lorentz internal indeces; we shall also choose the spacetime manifold to be without boundary in order to avoid in the induced action the presence of surface terms which are not relevant to our following discussion.

Let us assume that matter is described on a microscopic (or submicroscopic) level by one or more fundamental, massless, canonical, Dirac field $\psi(x)$. For the sake of simplicity we shall consider here a single field, the generalization to the case of n replicas of ψ is straight-forward. In order to define a spinor field like ψ on a manifold U which is not simply Minkowsky spacetime we must erect at every point an orthonormal frame where spinorial representations of the Lorentz group are allowed. This procedure is equivalent to define a global, non degenerate vierbein field $e^a{}_\mu(x)$ on U. Then an orthonormal frame $\hat\theta^a(x)$ is obtained from the coordinate basis through the relation $\hat\theta^a(x) = e^a{}_\mu(x)\, dx^\mu$. The matter field transforms under a local Lorentz rotation of $\hat\theta^a$ according to:

$$\psi'(x) = e^{-\frac{i}{2}\omega_{ab}(x)\Sigma^{ab}}\psi(x) \quad , \tag{1}$$

where $\Sigma^{ab} \equiv \frac{i}{4}[\gamma^a, \gamma^b]$ and γ^a are the constant Dirac matrices satisfying

$$\{\gamma^a, \gamma^b\} = -2\eta^{ab} = -2\,\mathrm{diag}(-1,+1,+1,+1) \quad . \tag{2}$$

Local symmetries require the introducing of appropriate covariant derivatives into the matter Lagrangian. General covariance is already guaranteed by the vierbein $e^a{}_\mu$, so a Lorentz connection $\tilde B_\mu(x) = \tilde B(x)_{ab\mu}\Sigma^{ab}$ is enough to build up a total covariant derivative

$$D_a \equiv e_a{}^\mu D_\mu = e_a{}^\mu(\partial_\mu - \frac{i}{2}\tilde B(x)_{ab\mu}\Sigma^{ab}) \quad . \tag{3}$$

To get an insight into the geometry of the SO(3,1) gauge theory, we have considered up to this point, let us look at the commutator of the covariant derivative (3). In a coordinate basis we have

$$[D_\mu, D_\nu] = -\frac{i}{2}F_{ab\mu\nu}\Sigma^{ab} \tag{4}$$

$$F_{ab\mu\nu} = \partial_\mu \tilde B_{ab\nu} - \partial_\nu \tilde B_{ab\mu} + \tilde B_{ac\mu}\tilde B^c{}_{b\nu} - \tilde B_{ac\nu}\tilde B^c{}_{b\mu} \quad . \tag{5}$$

Geometry manifestly comes into play if we use the metricity property of $e^a{}_\mu$ to trade the spin connection $\tilde B_{ab\mu}$ for a coordinate basis affine connection $\tilde\Gamma^\rho{}_{\lambda\mu}$[15)]:

$$\partial_\lambda e^a{}_\mu + \tilde B^a{}_{b\lambda}e^b{}_\mu - \tilde\Gamma^\nu{}_{\lambda\mu}e^a{}_\nu = 0 \quad . \tag{6}$$

In General Relativity one assumes $\tilde\Gamma^\nu{}_{\lambda\mu}$ to be the Levi-Civita connection, i.e. a metrical torsion free connection . This statement sounds at least unnatural in a spinor pregeometric theory where:

i) the metricity condition (6) is not a constraint but derives from the definitions of the spinor , Lorentz and affine connections[6)];

ii) the spin current density is expected to produce a torsion field, i.e. a non-symmetric part in $\tilde\Gamma^\nu{}_{\lambda\mu}$, in a close analogy with curvature field generated by the energy momentum density [5].

Motivated by the previous arguments we choose a coordinate basis affine connection of the form [*]

$$\tilde{\Gamma}^{\nu}{}_{\lambda\mu} = \left\{ {}^{\nu}_{\lambda\mu} \right\} + K^{\nu}{}_{\lambda\mu} \quad , \tag{7}$$

where $\left\{ {}^{\nu}_{\lambda\mu} \right\}$ is the Christoffel symbol,

$$K^{\nu}{}_{\lambda\mu} = T^{\nu}{}_{\lambda\mu} + T_{\lambda\mu}{}^{\nu} + T_{\mu\lambda}{}^{\nu} \tag{8}$$

is the contortion tensor and finally

$$T^{\nu}{}_{\lambda\mu} = \frac{1}{2}\left(\tilde{\Gamma}^{\nu}{}_{\lambda\mu} - \tilde{\Gamma}^{\nu}{}_{\mu\lambda} \right) \quad . \tag{9}$$

is the torsion tensor. Then we can relate the Lorentz field strength with the Riemann-Cartan curvature tensor by

$$F^{a}{}_{b\mu\nu} = e^{a}{}_{\alpha}\, e_{b}^{\beta}\, \tilde{R}^{\alpha}{}_{\beta\mu\nu} \tag{10}$$

$$\tilde{R}^{\alpha}{}_{\beta\mu\nu} = \partial_{\mu}\tilde{\Gamma}^{\alpha}{}_{\beta\nu} - \partial_{\nu}\tilde{\Gamma}^{\alpha}{}_{\beta\mu} + \tilde{\Gamma}^{\alpha}{}_{\mu\gamma}\tilde{\Gamma}^{\gamma}{}_{\nu\beta} - \tilde{\Gamma}^{\alpha}{}_{\nu\gamma}\tilde{\Gamma}^{\gamma}{}_{\mu\beta} \quad . \tag{11}$$

From the foregoing discussion we infer that in a spinor pregeometric model the underlying classical geometry is of the Riemann-Cartan type. In the following part of this note we shall assign a dynamical role to this background geometry through the vacuum fluctuations of the matter field

(*) From now on we shall use a tilde to indicate a geometric quantity in U^4. Purely Riemannian objects are written without tilde.

Let us start from the classical Dirac Lagrangian

$$L_D = \frac{i}{2}\bar{\Psi} e_a{}^{\mu}\gamma^a \left(\partial_\mu - \frac{i}{2}\tilde{B}_\mu\right)\Psi - \frac{i}{2}\bar{\Psi}\left(\overleftarrow{\partial}_\mu + \frac{i}{2}B_\mu\right) e_a{}^{\mu}\gamma^a\Psi \quad . \tag{12}$$

By splitting $\tilde{B}_{ab\mu}$ into the sum of the Ricci Rotation Coefficients $\omega_{ab\mu}(e)$ and a contortion part $K_{ab\mu}$, according to

$$\tilde{B}_{ab\mu} = \omega_{ab\mu}(e) + K_{ab\mu} \tag{13}$$

and taking into account the anti-commutation relation

$$\left\{ \gamma^a, \Sigma^{bc} \right\} = -\epsilon^{abcd}\gamma^5\gamma^d \quad , \tag{14}$$

where ϵ^{abcd} is the totally anti-symmetric tensor in 4-dimensions and

$$\gamma^5 = i\gamma^0\gamma^1\gamma^2\gamma^3 \tag{15}$$

one can get L_D in the form

$$L_D = \frac{i}{2}\bar{\Psi}\gamma^\mu\left(\partial_\mu - \frac{i}{2}\omega_{ab\mu}\Sigma^{ab} + \frac{3}{4}i\gamma^5 A_\mu\right)\Psi - \frac{i}{2}\bar{\Psi}\left(\overleftarrow{\partial}_\mu + \frac{i}{2}\omega_{ab\mu}\Sigma^{ab} + \right.$$
$$\left. + \frac{3}{4}i\gamma^5 A_\mu\right)\gamma^\mu\Psi \quad \equiv \frac{i}{2}\bar{\Psi}(\not{D}\Psi) - \frac{i}{2}(\bar{\Psi}\overleftarrow{\not{D}})\Psi \tag{16}$$

A_μ is the dual of the torsion, i.e.

$$A_\mu = \frac{1}{3!}\epsilon_{\mu\alpha\beta\gamma}T^{\alpha\beta\gamma} \tag{17}$$

The Dirac Lagrangian (16) manifestly shows the peculiar form of the spin-torsion coupling introducing a U(1) axial gauge symmetry into the model. We remark that this is a direct consequence of the general covariance and local Lorentz symmetry we have required at the beginning. Furthermore, we notice that no dimensional coupling constant is contained in L_D , hence the theory is not only chiral invariant but also Weyl invariant[11], i.e. we can freely rescale the field according to

$$\begin{aligned} e^a{}_\mu &\longrightarrow e^{\Omega(x)}e^a{}_\mu \\ e_a{}^\mu &\longrightarrow e^{-\Omega(x)}e_a{}^\mu \\ \psi(x) &\longrightarrow e^{-\frac{3}{2}\Omega(x)}\psi(x) \\ \tilde{B}_{ab\mu} &\longrightarrow \tilde{B}_{ab\mu} \end{aligned} \tag{18}$$

Now, it is clear that in order to obtain General Relativity characterized by the dimensional Newton constant we have to break this local scale invariance. The breaking of the Weyl invariance can be obtained by introducing an ultraviolet cut-off at the quantum level. Then all the dimensional constants in the low energy induced theory result to be fixed by the quantum cut-off acting as a fundamental mass scale.

At the classical level the $e^a{}_\mu(x)$ and $A_\nu(x)$ are non-dynamical in the sense that no kinetic terms appear for them in (16). The induced theory approach consists in the assignement to the regularized part of the quantum effective action the role of kinetic term for the background fields.

The generating functional, describing the quantum properties of the model is

$$Z[e^a{}_\mu(x), A_\nu(x)] = \int d\mu[\bar{\Psi},\Psi]\, e^{i\int dV_x L_D} \equiv e^{i\Gamma(e^a{}_\mu, A_\nu)} \, . \tag{19}$$

The n-points Green functions can be obtained by coupling the matter fields Ψ and $\bar{\Psi}$ to some external, fermionic, sources, say η and $\bar{\eta}$, and then functionally differentiating with respect to the sources. Here we are not concerned with vacuum expectation values field operators, so we shall omit $\bar{\eta}\Psi$ and $\bar{\Psi}\eta$ terms in the functional integral (19). Rather we are interested in the effective action Γ as a functional of the background fields $e^a{}_\mu(x)$, $A_\nu(x)$. L_D is bi-linear in Ψ and $\bar{\Psi}$ so we can easily sum over the spinor fields according to the functional integration rule over anti-commutating variables; the formal result of this operation is:

$$\Gamma = -i \ln\det(i e_a{}^\mu \gamma^a \mathcal{D}_\mu) = -\frac{i}{2}\ln\det\Delta \, . \tag{20}$$

The second order Dirac operator $\Delta \equiv (i\not{\mathcal{D}})(-i\not{\mathcal{D}})^\dagger$ reads

$$\begin{aligned} \Delta &= g^{\mu\nu}\nabla_\mu \mathcal{D}_\nu + X(x) \\ &= g^{\mu\nu}(\mathcal{D}_\mu - \Gamma_{\mu\nu}{}^\lambda)\mathcal{D}_\nu + i\,\Sigma^{\mu\nu}Y_{\mu\nu} \end{aligned} \, , \tag{21}$$

where

$$\mathcal{D}_\mu \equiv \partial_\mu - \frac{i}{2}\omega_{ab\mu}\Sigma^{ab} - \frac{3}{4}i\gamma^5 A_\mu \tag{22}$$

$$Y_{\mu\nu} \equiv [\mathcal{D}_\mu, \mathcal{D}_\nu] = -\frac{i}{2}R_{ab\mu\nu}\Sigma^{ab} - \frac{3}{4}i\gamma^5 F^5_{\mu\nu} \tag{23}$$

$$R_{ab\mu\nu} = \partial_\mu \omega_{ab\nu} - \partial_\nu \omega_{ab\mu} + \omega_{ac\mu} \omega^c{}_{b\nu} - \omega_{ac\nu} \omega^c{}_{b\mu} \tag{24}$$

$$F^5_{\mu\nu} = \partial_\mu A_\nu - \partial_\nu A_\mu \tag{25}$$

We now give the proper-time representation of the effective action, i.e.

$$\Gamma = \frac{i}{2} \int_0^\infty \frac{ds}{s} \mathrm{Tr}\, e^{-is\Delta} = \frac{i}{2} \int_0^\infty \frac{ds}{s} \mathrm{Tr}\, K(x,x;s) \tag{26}$$

and then select the divergent part of Γ by substituting in (26) the Heat-Kernel $K(x,x,s)$ with its asymptotic expansion

$$\mathrm{Tr}\, K(x,x,s) = i\,(4\pi i s)^{-2} \sum_0^\infty {}_m\, a_m(x)\,(is)^m \tag{27}$$

The first three terms in (27) are divergent when integrated over s, and must be regularized by introducing an ultraviolet cut-off Λ and an infrared one, say μ. After this procedure we find for the induced action

$$\Gamma^{IND} = \frac{1}{2(4\pi)^2} \int dV_x \left(\frac{1}{2} \Lambda^4 a_0 + \Lambda^2 a_1 + a_2 \ln \frac{\Lambda^2}{\mu^2} \right) , \tag{28}$$

where[16]

$$a_0 = \mathrm{Tr}\, 1\!\!1 = 4$$

$$a_1 = \mathrm{Tr}\left(1\!\!1 \frac{R}{6} - X \right) = -\frac{1}{3} R$$

$$a_2 = \mathrm{Tr}\Big(\frac{1}{12} Y_{\mu\nu} Y^{\mu\nu} + \frac{1}{2} X^2 - \frac{R}{6} X + \frac{1}{6} \mathcal{D}^2 X + \\ - \frac{1}{30} \Box R + \text{curvature square terms} \Big) = \frac{3}{8} (F^5_{\mu\nu})^2 + \dots \tag{29}$$

The dots in the final expression for a_2 correspond to quadratic curvature terms which are not relevant for the following discussion and therefore we shall not take into account. So we obtain for the induced action

$$\Gamma^{IND} = \int dV_x \left[- \frac{\Delta^4}{(4\pi)^2} + \frac{\Delta^2}{(4\pi)^2} \frac{R}{6} - \frac{3}{16(4\pi)^2} \ln \frac{\Delta^2}{\mu^2} \cdot (F^5_{\mu\nu})^2 \right] . \tag{30}$$

By matching (30) with the classical action

$$I^{CL} = \int dV_x \left[\frac{1}{16\pi G_N} (R - 2\lambda) - \frac{1}{4g^2} (F^5_{\mu\nu})^2 \right] \tag{31}$$

we identify the induced coupling constants

$$\frac{\lambda}{8\pi G_N} = \frac{\Lambda^4}{(4\pi)^2} \tag{32}$$

$$\frac{1}{16\pi G_N} = \frac{\Lambda^2}{6(4\pi)^2} \tag{33}$$

$$\frac{1}{4g^2} = \frac{3}{16(4\pi)^2} \ln \frac{\Lambda^2}{\mu^2} . \tag{34}$$

In order to give G_N the correct value we have to choose Λ near the Planck mass. The drawback of this choice is a huge induced cosmological constant which can be fine-tuned to zero only by adding a suitable , constant counterterm to the initial matter action. This is an unsatisfactory feature that Induced Quantum Gravity shares, as far as we know, with all the presently available quantum field theories coupled to gravity. Apart from some suggestions by Hawking[17] there does not seem to be any convincing argument to justify theoretically the present vanishingly small value of the cosmological constant. Even supersymmetry, once it is broken, produces an enormous value for λ .

As far as torsion is concerned we notice that the induced term describes the propagation of a massless axial vector field, so in our model torsion appears to be a long-range dynamical field. One could wonder why the cut-off does not simultaneously break the Weyl and the chiral invariance. We think that the reason is two-fold. Firstly the two symmetries act in different "worlds", in the sense that while local rescaling refer to spacetime, axial gauge transformations correspond to chiral rotations in the internal symmetry space. Secondly, our computational scheme, and in particular the proper time regularization method, is gauge invariant. Therefore the unbroken U(1) axial symmetry preserves A_μ from acquiring a mass, which otherwise should be of the order of the Planck mass. Unfortunately we have not a similar mechanism to suppress the undesirable cosmological constant.

We recall that in the classical Einstein-Cartan-Sciama-Kibble theory[18], torsion does not propagate in vacuum. It is "frozen" inside the spin density distribution. Such a behaviour causes severe patologies on a quantum level. A possible way out is to modify the initial Lagrangian inserting "by hand" suitable kinetic terms[13),19]. But then one inserts into the theory new unspecified coupling constants, one for each new term. Moreover the particle spectrum, in the linearized theory, results to be physically acceptable only if some suitable constraints are imposed on these new coupling constants.

A different proposal suggests that torsion could propagate through fermion loops at the quantum level. Despite the appearence, this approach is basically different from the one we presented here in the sense that torsion squared terms are already present in the classical action one starts from. In our model all the curvature and torsion terms arise only at the quantum level actually we started from the Dirac Lagrangian alone. Let us conclude with some comments concerning the possible quantum behaviour of the induced action (30). The Newton constant (33) and gauge charge (34) have the correct positive sign. They are related by an $\alpha - G_N$ type[20] relation of the form

$$\frac{1}{4g^2} = \frac{3}{16(4\pi)^2} \ln \frac{6\pi}{\mu^2 G_N} \qquad (35)$$

This result suggests that we can reverse the argument we used to fix the value of Λ, and say that gravity provides a universal cut-off $\Lambda \sim G_N^{-1/2}$. This conclusion has also been obtained some time ago by non-perturbative summations of particular types of Feynman graphs in quantum gravity[21].

The induced action, including the quadratic curvature terms we discarded in $\mathcal{a}_2(x)$, is perturbatively renormalizable both in the gravitational[22] and the torsion sector. The second part of the statement being guaranteed by the Maxwell like form of the Lagrangian for A_μ.

Finally, if the gravitational vacuum is modelled as a super conducting medium, we could expect in a more sophysticated version of our naive pregeometric model, the A_μ field to be confined inside thin flux tubes. This sort of "Meissner effect" has recently been proposed to explain the absence of torsion at large distances as a dynamical phenomenon[23].

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