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LETTER TO THE EDITOR

On the first-order formalism in quantum gravity

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Abstract. We show that quantum gravity in the first-order form is equivalent to the quantised Einstein theory only in the phases where $\langle e^a_\mu \rangle = e_{0\mu}^{\ a}$, $\det e_0 \neq 0$. If $\det e_0 = 0$ one must take into account configurations with non-zero torsion. We also discuss the analogy with the first-order formulation for gauge theories and the connection-dependent counterterms for the spinor theory.

Here we shall make several remarks about the quantisation of gravity in the first-order form (or the Einstein–Cartan theory, see e.g. Hehl *et al* (1976))

$$\begin{aligned} &\mathscr{L}_{\mathrm{E}} = -k^{-2} e^a_\mu \, {}^*R^{*\mu\nu}_{\ ab} \, e^b_\nu = -k^{-2} R^{ab}_{\mu\nu} \, {}^*\pi^{*\mu\nu}_{\ ab}, \\ &I = \int \mathscr{L}_{\mathrm{E}} \, \mathrm{d}^4 x, \qquad \pi^{ab}_{\mu\nu} = e^a_{[\mu} e^b_{\nu]}, \qquad\qquad (1) \\ &{}^*R^{*\mu\nu}_{\ ab} = \tfrac{1}{4} \varepsilon_{abcd} \varepsilon^{\mu\nu\lambda\rho} R^{cd}_{\lambda\rho}(\omega), \end{aligned}$$

where e^a_μ and ω^{ab}_μ are the vierbein and the SO_4 connection (we use the Euclidean formulation). Several arguments support the choice of (1) as a starting point for quantisation: (i) the (e, ω) formulation is necessary from the point of view of the gauge approach to gravity and supergravity (Hehl *et al* 1976, Ne'eman 1978, Ne'eman and Regge 1978); (ii) it is the vierbein and not the metric that should be used in the path integral in order to avoid the problem of maintaining the correct signature of the metric (see e.g. Hawking 1979); (iii) working in the first-order formalism, one is able to write down not only the gravitational (1) but also the *matter Lagrangian* in completely *polynomial form*

$$\begin{aligned} \mathscr{L}_{\mathrm{m}} = {}& [\mathrm{i}\phi^a e_{3a}^{\ \mu} \, \partial_\mu \varphi + (\tfrac{1}{2}\phi_a^2 + V(\varphi) e_4)] + [\tfrac{1}{4}(F^\alpha_{ab})^2 e_4 + (\mathrm{i}/2g) F^{ab}_\alpha \, {}^*\pi^{*\mu\nu}_{\ ab} \mathscr{F}^\alpha_{\mu\nu}(A)] \\ & + [\tfrac{1}{2}\mathrm{i}\bar\psi\gamma^a \overleftrightarrow{\mathscr{D}}_\mu(\omega)\psi e_3^{\ \mu} + m\bar\psi\psi e_4] + [-\tfrac{1}{2}\bar\psi_\mu \gamma_a \gamma_5 \mathscr{D}_\nu(\omega)\psi_\rho \varepsilon^{\mu\nu\rho\sigma} e^a_\sigma], \qquad (2) \\ e_{3a}^{\ \mu} = {}& (1/3!)\varepsilon^{\mu\nu\gamma\rho} \varepsilon_{abcd} e^b_\nu e^c_\lambda e^d_\rho = e^\mu_a e_4, \qquad \mathscr{F}^\alpha_{\mu\nu} = \partial_\mu A^\alpha_\nu - \ldots; \end{aligned}$$

ϕ and F are the natural auxiliary fields for the scalar and gauge fields (cf Taylor 1978) and 'i's are necessary for the correspondence with the Euclidean path integral

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quantisation. (For example,

$$\int \mathrm{d}A_\mu \exp\left[-\frac{1}{4g^2}\mathscr{F}^2_{\mu\nu}(A)\right]$$

$$= \int \mathrm{d}G_{\mu\nu}\,\mathrm{d}A_\rho\,\delta[G_{\mu\nu}-(1/g)\mathscr{F}_{\mu\nu}]\exp(-\tfrac{1}{4}G^2_{\lambda\rho})$$

$$= \int \mathrm{d}F_{\mu\nu}\,\mathrm{d}G_{\lambda\rho}\,\mathrm{d}A_\sigma \exp\{\tfrac{1}{2}\mathrm{i}F^{\mu\nu}[G_{\mu\nu}-(1/g)\mathscr{F}_{\mu\nu}]-\tfrac{1}{4}G^2_{\lambda\rho}\}$$

$$= \int \mathrm{d}F_{\mu\nu}\,\mathrm{d}A_\lambda \exp[-\tfrac{1}{4}F^2_{\mu\nu}-(\mathrm{i}/2g)F^{\mu\nu}\mathscr{F}_{\mu\nu}(A)].)$$

It is worth noting that matter Lagrangians arise precisely in the first-order form in the 'group manifold' approach to the construction of extended supergravities (D'Auria *et al* 1980, Fre 1981).

This polynomiality seems very important for non-perturbative path integral quantisation, either in a continuum limit (Taylor 1979, Nouri-Moghadam and Taylor 1980) or on a lattice (Smoline 1979, Das *et al* 1979, Mannion and Taylor 1981). The absence of a polynomial 'free' term for e or ω (excluding the trivial Λ term) sharpens the problem of phases in the theory. It seem especially natural in the first-order formalism to assign the following dimensions to the fields (De Alfaro *et al* 1980a, b):

$$[e^a_\mu, \omega^{ab}_\mu, A^\alpha_\mu, \psi_\mu] = cm^{-1}, \qquad [\varphi, \phi_a, F^\alpha_{ab}, \psi] = cm^0. \tag{3}$$

Now k in (1) is dimensionless while in the 'Einstein phase'

$$\langle e^a_\mu\rangle = \varkappa^{-1}\delta^a_\mu, \qquad \varkappa^2 = 16\pi G. \tag{4}$$

Here we want to point out that the choice (4) (necessary in the second-order formalism to provide a meaning for $g^{\mu\nu}$) is not a unique phase in the first-order theory (1).

Let us first consider the classical field equations for (1),

$$R^{ab}\wedge e^c\varepsilon_{abcd} = 0, \qquad \text{or } {}^*R^{*\mu\nu}_{ab}e^b_\nu = 0, \tag{5}$$

$$T^a\wedge e^b\varepsilon_{abcd} = 0, \qquad \text{or } \mathscr{D}_\mu{}^*\pi^{*\mu\nu}_{ab} = 0, \tag{6}$$

where $T^a = \mathscr{D}e^a$ is the torsion two-form. It is important to realise that equations (5) and (6) are equivalent to the Einstein equations *only when* $\det e \neq 0$ (then (6) implies $T^a = 0$ or $\omega = \omega_0(e)$). In fact, we found the following solutions of (5) and (6) *with* $\det e = 0$ *and non-zero torsion* for the SO_4 and SO_3 (static)-symmetrical cases (here we assume that the Λ term is included in (5)):

$$\begin{aligned}
&SO_4\colon e^a = \alpha\,\mathrm{d}n^a, \qquad n^a = x^a/\rho, \qquad \rho^2 = x_a x^a,\\
&\omega^{ab} = \varphi_1 n^{[a}\,\mathrm{d}n^{b]} + \varphi_2\varepsilon^{abcd}n_c\,\mathrm{d}n_d,\\
&\alpha, \varphi_1, \varphi_2 = \text{constant}, \qquad \varphi_1 - \tfrac{1}{4}\varphi_1^2 - \varphi_2^2 = \tfrac{1}{3}\Lambda\alpha,\\
&T^a = \alpha\varphi_2\varepsilon^{abcd}n_b\,\mathrm{d}n_c\wedge\mathrm{d}n_d \neq 0, \qquad R^{ab}\neq 0.
\end{aligned}\tag{7}$$

$$\begin{aligned}
&SO_3\colon e^0 = f\,\mathrm{d}t, \qquad e^i = \alpha\,\mathrm{d}n^i, \qquad n^i = x^i/r, \qquad r^2 = x^i x^i,\\
&\omega^{ij} = \varepsilon^{ijk}\psi n^k\,\mathrm{d}t, \qquad \omega^{0i} = \varphi n^i\,\mathrm{d}t,\\
&\varphi = -\tfrac{1}{2}\Lambda\alpha f, \qquad \alpha, f = \text{constant}, \qquad \psi(r) = \text{arbitrary},\\
&T_a \neq 0, \qquad R_{ab}\neq 0.
\end{aligned}\tag{8}$$

One can mention several properties of these solutions: (a) torsion and curvature swiftly decrease with ρ (or r)$\to\infty$ (cf the proposals of torsion instantons or monopoles by Hanson and Regge (1978)); note also that (7)

$$(e^a_\mu = (\alpha/\rho)(\delta^a_\mu - x^a x_\mu/\rho^2), \qquad \omega^{ab}_\mu = 2Ax^{[a}\delta^{b]}_\mu + B\varepsilon^{ab}_{\mu\nu}x^\nu,$$
$$A = \varphi_1/2\rho^2, \qquad B = -\varphi_2/2\rho^2)$$

can be compared with the one-instanton

$$(e^a_\mu = (6/\Lambda)^{1/2}[a/(\rho^2+a^2)]\delta^a_\mu, \qquad A = 2/(\rho^2+a^2), \qquad B = 0)$$

and meron ($e^a_\mu \sim (1/\rho)\delta^a_\mu$) solutions (see e.g. De Alfaro *et al* 1980a); (b) (7) and (8) have a universal scale dependence on x^μ, naturally corresponding to the choice (3); (c) the metrics in (7) and (8) measure distances only in the directions orthogonal to n^a or n^i. Hence one can speak about space–times with $\det e = 0$ as having 'metrical dimension' $d < 4$ (and so (7), (8) or their multidimensional generalisations may have some relation to a mechanism of dimensional reduction and also to the effect of quantum fluctuations of the space–time dimension (Banks 1980)). Observe also that the behaviour of matter on the background of (7) or (8) also corresponds to the 'dimensional reduction' (e.g. the corresponding solution for φ in (2) is ρ or r independent).

The main question is why solutions with $\det e = 0$ are to be considered as physical at all. We think that they must be included in the path integral (Hawking (1979) pointed out that changes of the space–time topology occur through these degenerate metrics). Thus one can conjecture the possibility of the 'non-Einsteinian' phases $\langle e\rangle = e_0$, $\det e_0 = 0$ in the theory (for the discussion of the phases different from (4) see also Taylor (1979), Nouri-Moghadam and Taylor (1980), De Alfaro *et al* (1980a, b)). The simplest solution of (5), (6) with $\det e = 0$ is $e^a_\mu = 0$, ω^{ab}_μ = arbitrary. The corresponding phase $\langle e\rangle = 0$ is essentially 'non-classical': in some sense the classical metrical space–time does not exist in it (cf Taylor 1979). Comparing $\langle e^a_\mu\rangle = \varkappa^{-1}\delta^a_\mu$ and $\langle e^a_\mu\rangle = 0$ with the φ^4 theory phases $\langle\varphi\rangle = a$ and $\langle\varphi\rangle = 0$, one can speculate that the $\langle e\rangle = 0$ phase is restored at high temperatures and curvatures (initial singularity or the final state of collapse).

Now let us consider the formal Euclidean path integral for (1),

$$Z = \int \mathrm{d}e\, \mathrm{d}\omega\, \exp(k^{-2}e^*R^*e + et + \omega S), \tag{9}$$

where t and S are the sources (energy–momentum and spin) and we imply some boundary conditions for e and ω (or the choice of the vacuum). First we shall discuss the integration over ω, assuming $\langle\omega\rangle = 0$. Adding the square of torsion to (1), one has

$$\begin{aligned}\mathscr{L}'_{\mathrm{E}} &= (\sqrt{g}/k^2)[-R(\omega) + \tfrac{1}{2}\alpha(T^a_{\mu\nu})^2]\\ &= \frac{\sqrt{g}}{k^2}\Big(-R[\omega_0(e)] + \frac{\alpha-1}{2}(\bar\omega - \bar\omega_0)^2\\ &\quad + \frac{\alpha+2}{3}(\check\omega - \check\omega_0)^2 + 6(2\alpha+1)(\hat\omega - \hat\omega_0)^2\Big),\end{aligned} \tag{10}$$

where

$$\omega_{ab\mu}e^\mu_c = \bar\omega_{abc} + \tfrac{2}{3}\delta_{c[b}\check\omega_{a]} + \varepsilon_{abcd}\hat\omega^d, \qquad \bar\omega_{[abc]} = 0, \qquad \check\omega_a = \omega^{\ b}_{ab}.$$

So the action (1) ($\alpha = 0$) is non-positive not only in the 'e sector' (Hawking 1979) but also in the 'ω sector' (cf Deser and Nicolai 1981). Thus the integral over ω in (9) should be defined either by $\bar\omega \to \mathrm{i}\bar\omega$ or by going to (10) with $\alpha \geq 1$ (this does not disturb the equivalence with the Einstein theory when $\det e \neq 0$ and $S=0$ but makes the action non-polynomial). The result of ω integration (trivial from (10)) can also be written (for $S=0$) in the form (see (1))

$$Z=\int \frac{\mathrm{d}e}{(\det K)^{1/2}} \exp\left(-\frac{1}{k^2}(\partial_\mu {}^*\pi^{*\mu\nu}_{\ \ \alpha})K^{-1\alpha\beta}_{\ \ \ \gamma\lambda}(\partial_\rho {}^*\pi^{*\rho\lambda}_{\ \ \beta})+et\right), \tag{11}$$

where $K^{\mu\nu}_{\beta\gamma}=f^\alpha_{\beta\gamma}{}^*\pi^{*\mu\nu}_{\ \ \alpha}\gtrless 0$ and so the analytic continuation is again assumed; $\alpha=[ab]$ and $f^\alpha_{\beta\gamma}$ here are the SO_4 structure constants. The first term in (11) is the Einstein action (with the boundary term), and so e^a_μ contribute in it only through $\pi^{ab}_{\mu\nu}$ and its non-polynomiality stems from K^{-1}. It should be stressed that (11) is valid only when $\det K\neq 0$ or $\det e\neq 0$, and therefore (9) is equivalent to the Einstein theory only up to the contribution of configurations with $\det e=0$. Hence the consequence of (9) is (11) *plus* the result of ω integration for the case of $\det e=0$.

Next we shall remark on the connection of (1), (9), (11) with the first-order (and field strength) formulation of the SO_4 gauge theory in the flat space–time (see e.g. Halpern 1979, Seo and Okawa 1980) (cf (2))

$$\mathcal{L}=\tfrac{1}{4}(F^\alpha_{\mu\nu})^2+(\mathrm{i}/2g)F^{\mu\nu}_\alpha\mathcal{F}^\alpha_{\mu\nu}(A). \tag{12}$$

Assuming that (i) one can neglect the F^2 term in (12) and (ii) $F^{\mu\nu}_\alpha$ is the *composite* field in terms of some new one e^a_μ, $F^{\mu\nu}_{ab}=2\mathrm{i}{}^*\pi^{*\mu\nu}_{\ \ ab}$, we conclude that (12) coincides with (1) with $g=k^2$. Moreover, integrating over A_μ in the path integral for (12), we obtain

$$(\mathcal{K}^{\mu\nu}_{\beta\gamma}=f^\alpha_{\beta\gamma}F^{\mu\nu}_\alpha, \qquad \mathcal{F}^\alpha_{\mu\nu}\equiv\partial_\mu A^\alpha_\nu-\partial_\nu A^\alpha_\mu+f^\alpha_{\beta\gamma}A^\beta_\mu A^\gamma_\nu)$$

$$Z=\int\frac{\mathrm{d}F}{(\det\mathcal{K})^{1/2}}\exp\left(-\tfrac{1}{4}F^2+\frac{\mathrm{i}}{2g}(\partial_\mu F^{\mu\nu}_\alpha)\mathcal{K}^{-1\alpha\beta}_{\ \ \ \nu\lambda}(\partial_\rho F^{\rho\lambda}_\beta)\right). \tag{13}$$

The comparison of (11) and (13) shows that the Einstein Lagrangian is exactly the '$\partial FF^{-1}\partial F$' part of the SO_4 gauge theory Lagrangian in the F formulation under the substitution $F\to 2\mathrm{i}{}^*\pi^*$.

A natural intention now is to integrate over e^a_μ in (9) in order to obtain the 'dual' to the Einstein formulation of the theory thus analogous to the ordinary A_μ (or ω) formulation of the gauge theory. It is important to realise that the gaussian integral over e (suitably defined in view of ${}^*R^*\gtrless 0$) is exactly calculable,

$$Z=\int\frac{\mathrm{d}\omega}{(\det {}^*R^*)^{1/2}}\exp(-\tfrac{1}{2}k^2 t\,{}^*R^{*-1}t+\omega S), \tag{14}$$

only in the $\langle e\rangle=0$ phase. Thus the theory (14) (possessing rather unusual properties like the absence of propagation of interaction of 't's in view of the algebraic nature of ${}^*R^*$† and the lack of the free ω term apart from $\mathrm{Tr}\log {}^*R^*\sim\delta^4(0)$) is *not* equivalent to

† To give the idea of the explicit form of ${}^*R^{*-1}$ we present here the expression for $\mathcal{K}^{-1}$ in (13) in the case of SU_2:

$$\mathcal{K}^{-1ij}_{\ \ \mu\nu}=[\varepsilon^{ijk}F^k_{\mu\nu}]^{-1}=(16/\det G)(-\tfrac{1}{4}G^{ij}N_{\mu\nu}+\tfrac{1}{8}G^{im}G^{jn}\overset{*}{F}{}^k_{\mu\nu}\varepsilon_{mnk}),$$
$$G^{ij}=F^i_{\mu\nu}\overset{*}{F}{}^j_{\mu\nu}, \qquad N_{\mu\nu}=\tfrac{1}{3}\varepsilon_{ijk}\overset{*}{F}{}^i_{\mu\lambda}F^j_{\lambda\rho}\overset{*}{F}{}^k_{\rho\nu}$$

(this general result seems to be absent in the literature, cf Halpern (1979), Seo and Okawa (1980) and references therein).

the Einstein theory in the usually assumed phase (4) (however, it *is* equivalent to (11) with the additional term mentioned above, but again only if $\langle e\rangle = 0$ is assumed in (11)). In the phases with $\langle e\rangle \neq 0$ the e integral in (9) cannot be exactly calculated (one must take into account the zero modes of $^*R^*$ or the solutions of (5), (6)). An obvious semiclassical approximation (in both $e = e_0 + h$ and $\omega = \tilde{\omega} + w$) gives (for $\det e_0 \neq 0$) the well known semiclassical result in the Einstein theory (*independently* of the order of the h and w integrations). If, however, $\det e_0 = 0$ one must take into account the configurations with non-zero torsion like (7), (8).

It should be noted that in the above formal discussion we have ignored (e.g. assuming a cut-off) possible counter-terms in (9) which may preclude an exact integration over e and ω. Observe also that the problem of gauge fixing does not arise in the separate e and ω integrations in (9) but only in the whole integral.

When matter is present we shall consider e and ω on an equal footing, calculating first the path integrals over the matter fields. As an example, let us discuss the divergences of $\log\det \hat{\mathscr{D}}(\omega)$ for the spin-$\frac{1}{2}$ theory in (2). This calculation is very simple (cf the straightforward but unnatural and unnecessarily complicated approaches of Goldthorpe (1980), Kimura (1981) and Nieh and Yau (1981)) if one observes that: (i) ω contributes to $\bar{\psi}\hat{\mathscr{D}}\psi$ only through the composite field B_μ,

$$\mathscr{L}_{1/2} = i\bar{\psi}\{\tfrac{1}{2}\hat{\mathscr{D}}[\omega_0(e)] + \gamma_5\gamma^\mu B_\mu\}\psi, \qquad B^\mu = \tfrac{1}{2}\varepsilon^{\mu\lambda\rho\sigma} e^a_\lambda T_{a\rho\sigma}; \tag{15}$$

(ii) $\tilde{B}_\mu = \gamma_5 B_\mu$ can be formally considered as an internal U_1 gauge field in view of $[\sigma_{ab}, \gamma_5] = 0$ or $[\omega_{0\mu}, \tilde{B}_\nu] = 0$, $\omega_{0\mu} = \frac{1}{2}\sigma_{ab}\omega^{ab}_{0\mu}$. Then, introducing $D_\mu = \partial_\mu + \omega_{0\mu} + \tilde{B}_\mu$, we obtain, using a well known algorithm (see e.g. Schwarz 1979), the corresponding counter-terms in the form

$$\begin{aligned} \Delta\mathscr{L}_{1/2} &= -(n-4)^{-1} b_4, \\ (4\pi)^2 b_4 &= -\tfrac{2}{3} W^2_{\mu\nu} - \tfrac{7}{360} R^* R^* - \tfrac{1}{10}(R^2_{\mu\nu} - \tfrac{1}{3}R^2) - \tfrac{1}{30}\mathscr{D}^2 R, \end{aligned} \tag{16}$$

where $R_{\dots} = R_{\dots}(e)$, $W_{\mu\nu} = \partial_\mu B_\nu - \partial_\nu B_\mu$. Thus the theory (1) with spinors is non-renormalisable not only in the 'e sector' but also in the 'ω sector' (one needs the bare $W^2_{\mu\nu}$ term leading to the 'propagating torsion'). Observe also that the interaction of ψ and ω through $\hat{\mathscr{D}}(\omega)$ is *not* the interaction of fermions with the internal SO_4 gauge field because $[\gamma_a, \sigma_{cd}] \neq 0$. This is the cause of the appearance of $W^2_{\mu\nu}$ instead of the 'ordinary' $[R^{ab}_{\mu\nu}(\omega)]^2$ term in (16). Therefore, various $R^2(\omega)$ additions to (1) (see e.g. Neville 1978, Sezgin and van Nieuwenhuizen 1980, Christensen 1980) seem to be unnatural from the point of view of the interaction of ω with the basic spin-$\frac{1}{2}$ matter field. However, the generalisation of (16) for the spin-$\frac{3}{2}$ field will be more complicated because all irreducible parts of the connection contribute to the gravitino Lagrangian and thus in the ψ_μ–ω interaction (note that the formal integration over ψ_μ before the integration over ω was recently considered by Deser and Nicolai (1981)).

In conclusion, let us point out that if it is possible to construct finite supergravity theories strictly in the first-order form (with ω being independent off-shell) then not only e- but also ω-dependent divergences will mutually cancel. As a result, one will be able to justify the formal integration over ω and (or) the use of the polynomiality of the bare action in the study of phases of the theory and in a non-perturbative calculation of the integral over e^a_μ.

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