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Composition Independence of the Possible Finite-Range Gravitational Force

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Abstract

A scalar–tensor theory of gravitation with scale invariance broken spontaneously is examined to show if and how the finite-range force added to the ordinary Newtonian force can be composition independent in the Eötvös experiment.

§(1): *Introduction*

Much progress has been and is going to be made in different types of experiments to test a possible finite-range gravitational force at macroscopic distances [1]. The potential between two massive objects may be given by [2]

$$V(r) = -G_{\infty}(m_1 m_2/r)\,(1 + ae^{-r/\lambda}) \tag{1}$$

One of the theoretical bases of (1) is provided by a version of the scalar-tensor theory of gravitation which offers a natural framework in formulating the concept of spontaneously broken scale invariance [3-6]; the ultimate aim of the approach is to understand the gravitational constant, on one hand, and the masses of elementary particles, on the other hand, in terms of a common origin. We proposed [3] a theory by introducing two scalar fields, one of which acquires a mass, giving rise to the second term of (1).

On the basis of this theory we now focus our attention on the Eötvös experiment. Although this experiment is not designed to measure how the gravitational force varies with distance, it should tell us whether the finite-range force is also composition independent. The answer to this question can be made affirmative from a theoretical point of view. We show this by improving our previous analysis to include two or more interacting matter fields.

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We emphasize that there is no *a priori* reason why the strength of the scalar fields should be proportional to the total energy of a gravitating system; unlike the tensor gravitational field with the total matter energy-momentum tensor as its source, scalar fields may have sources which are of quite a different nature. In our theory, fortunately, the couplings of the scalar fields are related intimately to the particle masses which are generated spontaneously in the originally massless theory. For this reason it is reasonable to expect that the scalar fields couple essentially to the trace of the *effective* energy-momentum tensor of the matter fields, some of which are now massive.

It should be kept in mind, however, that the present accuracy of the Eötvös experiment may not be sufficient to test detailed properties of the finite-range force in question; the additional contribution is proportional to the parameter a of (1), which could be smaller than 1, and to the ratio $(\lambda/R)^3$, which is $\lesssim 4 \times 10^{-9}$ for $\lambda \lesssim 10$ km, R being the radius of the earth. (The finite-range force gives no contribution in the experiments using the Sun as the gravity source [7].) It is premature to conclude whether composition independence has been established or not for the finite-range force, if any, before more detailed information is available for a and λ.

Nevertheless we assume for the moment that composition independence is honored by the finite-range force as well, and seek the condition under which the source of one of the scalar fields is proportional to the trace T of the *total* effective matter energy-momentum tensor $T_{\mu\nu}$. Although $-T = -T^{\mu}_{\ \mu}$ differs from T_{00} in any microscopic regions, they should give the same result if they are integrated over a *static macroscopic* object with negligible internal stress, like a rigid body. In order to demonstrate the required proportionality, we examine a simple nontrivial example of the interacting matter system of a Dirac field and a scalar field.

In Section 2 the basic Lagrangian is presented. We apply the technique of a conformal rescaling which is more convenient for the present purposes than to work with field equations as in Reference 3. The transformed Lagrangian is analyzed in detail in Section 3. In Section 4 the condition for the proportionality is obtained. The result is extended also to the system of many fermions and bosons. A simple choice of the coupling constants is suggested. The final Section 5 contains discussion and remarks on the nature of the theory.

§(2): *The Basic Lagrangian and Conformal Rescaling*

We assume the Lagrangian

$$\mathcal{L} = \mathcal{L}_g + \mathcal{L}_0 + \mathcal{L}_1 + \mathcal{L}_2 + \mathcal{L}_3 \tag{2a}$$

where

$$\mathcal{L}_g = b\left(\tfrac{1}{2} f^{-2} \phi^2 R - \tfrac{1}{2} \phi_{,\mu} \phi^{,\mu} - \tfrac{1}{2} \psi_{,\mu} \psi^{,\mu}\right) \tag{2b}$$

$$\mathcal{L}_0 = b(-\overline{\Psi}\not{D}\Psi - \tfrac{1}{2}\Phi_{,\mu}\Phi^{,\mu}) \tag{2c}$$

$$\mathcal{L}_1 = -b \sum_{r=0}^{4} c_r \phi^{4-r} \psi^r \tag{2d}$$

$$\mathcal{L}_2 = -b[(h_1\phi + h_2\psi)\overline{\Psi}\Psi + \tfrac{1}{2}(g_1^2\phi^2 + g_2^2\psi^2)\Phi^2] \tag{2e}$$

$$\mathcal{L}_3 = -be\overline{\Psi}\Psi\Phi \tag{2f}$$

where $b = \det(b^i{}_\mu)$ with $b^i{}_\mu$ the vierbein field; ϕ and ψ are the two "gravitational" scalar fields.[1] As matter fields we introduce a Dirac field Ψ and a real scalar field Φ which are coupled to each other through the Yukawa interaction $\mathcal{L}_3$.

The covariant derivative D_μ is given by

$$D_\mu\Psi = (\partial_\mu + i\,\tfrac{1}{4}A^{ij}{}_\mu\,\sigma_{ij})\,\Psi \tag{3a}$$

where the spinor connection $A_{ij\mu}$ is defined in the second-order formalism by

$$A_{ij\mu} = \tfrac{1}{2}b^k_\mu(c_{kij} - c_{ijk} - c_{jki}) - \tfrac{1}{2}\tau\phi^{-2}b^k_\mu\epsilon_{ijkl}A^l \tag{3b}$$

with

$$c_{kij} = (b^\mu_i b^\nu_j - b^\nu_i b^\mu_j)\,b_{k\mu,\nu} \tag{3c}$$

$$A^l = i\overline{\Psi}\gamma_5\gamma^l\Psi. \tag{3d}$$

The curvature tensor is defined by

$$R^{ij}{}_{\mu\nu} = -A^{ij}{}_{\mu,\nu} + A^i{}_{k\mu}A^{kj}{}_\nu - (\mu \longleftrightarrow \nu) \tag{4}$$

All the constants f, c_r, h's, g's, e, τ are real dimensionless so that the Lagrangian (2) possesses a complete scale invariance.

We apply a conformal rescaling

$$b^i_\mu(x) \longrightarrow \Lambda^{-1}(x)\,b^i_\mu(x) \tag{5a}$$

$$\Psi(x) \longrightarrow \Lambda^{1/2}(x)\,\Psi(x), \qquad \overline{\Psi}(x) \longrightarrow \Lambda^{1/2}(x)\,\overline{\Psi}(x) \tag{5b}$$

We easily find that

$$A_{ij\mu} \longrightarrow A_{ij\mu} - b_{i\mu}F_j + b_{j\mu}F_i \tag{6a}$$

with

$$F_i = b^\mu_i F_\mu = b^\mu_i\,(\log\Lambda)_{,\mu} \tag{6b}$$

[1] The Greek and Latin letters are used for world and Lorentz indices, respectively. The constant Dirac matrices γ_i obey $\tfrac{1}{2}\{\gamma_i, \gamma_j\} = \eta_{ij} = \mathrm{diag}\,(-+++)$. The Levi-Civita symbol ϵ^{ijkl} is normalized by $\epsilon^{0123} = +1$, while $\gamma_5 = i\gamma_0\gamma_1\gamma_2\gamma_3$. We use other notations as close as possible to those in Reference 3.

Substituting (6) into (4) we obtain

$$R \longrightarrow \Lambda^2 (R + 6\nabla_\mu F^\mu - 6F_\mu F^\mu) \tag{7}$$

Notice that the field Ψ does not appear explicitly here. The first term of (2b) can be brought to a pure Einstein-Hilbert action by choosing

$$\Lambda(x) = \kappa f^{-1} \phi(x) \tag{8}$$

The second term of (7) can be dropped because it finally results in a 4-divergence. The last term of (7) combines with the second term of (2b) to give

$$-(f^2/2\kappa^2)\, bZ^{-1} \phi^{-2} \phi_{,\mu}\phi^{,\mu} \tag{9a}$$

with

$$Z^{-1} = 1 + 6f^{-2} \tag{9b}$$

This term is shown to be an ordinary kinetic energy term of the field σ_1 defined by

$$\phi = v_1 \exp(\gamma\sigma_1) \tag{10a}$$

where

$$\gamma = \kappa f^{-1} Z^{1/2} \tag{10b}$$

Equation (10a) may be interpreted as meaning that ϕ has a "vacuum expectation value" v_1.

§(3): *The Transformed Lagrangian*

The remaining terms of (2) are also brought to simpler forms if we redefine the fields as

$$\psi \longrightarrow \kappa v_1 f^{-1} (v_2 + \sigma_2) \tag{11a}$$

$$\Psi \longrightarrow \kappa v_1 f^{-1} \Psi, \quad \Phi \longrightarrow \kappa v_1 f^{-1} \Phi \tag{11b}$$

The results will be discussed for various terms separately.

The first two terms of (2a) become

$$\mathcal{L}_g + \mathcal{L}_0 \longrightarrow b(1/2\kappa^2) R - b\tfrac{1}{2}\sigma_{1,\mu}\sigma_1{}^{,\mu} + be^{-2\gamma\sigma_1}\left[-\tfrac{1}{2}\sigma_{2,\mu}\sigma_2{}^{,\mu} - \overline{\Psi}\partial\!\!\!/\Psi - \tfrac{1}{2}\Phi_{,\mu}\Phi^{,\mu} + \tfrac{1}{8} c_{ijk}\,\epsilon^{ijkl} A_l - \tfrac{3}{4}\tau v_1^{-2} A_l A^l e^{-\gamma\sigma_1}\right] \tag{12}$$

The last two terms in the square bracket, related to the presence of the torsion, may not be important in most of macroscopic phenomena in which matter spins are always averaged.

In the potential term $\mathcal{L}_1$ we keep only terms at most quadratic in σ's:

$$b\mathcal{L}_1 \longrightarrow -bf^4\kappa^{-4}\left[F_0(\xi) + \kappa f^{-1} F_1(\xi)(-Z^{1/2}\xi\sigma_1 + \sigma_2) + \kappa^2 f^{-2} F_2(\xi) \cdot (\tfrac{1}{2}Z^{1/2}\xi\sigma_1 - \sigma_2) Z^{1/2}\sigma_1 + \kappa^2 f^{-2} F_3(\xi)\,\sigma_2{}^2 + O(\sigma^3)\right] \tag{13a}$$

where

$$\xi = \kappa f^{-1} v_2 \tag{13b}$$

$$F_0(\xi) = \sum_{r=0}^{4} c_r \xi^r \tag{13c}$$

$$F_1(\xi) = dF_0(\xi)/d\xi \tag{13d}$$

$$F_2(\xi) = \sum_{r=0}^{4} r^2 c_r \xi^{r-1} \tag{13e}$$

$$F_3(\xi) = \tfrac{1}{2} \xi^{-1} [F_2(\xi) - F_1(\xi)] \tag{13f}$$

The equivalence between (13) and the previous result given by equations (4)-(6b) in Reference 3 is obvious by replacing $x = v_2/v_1$ in Reference 3 by ξ.

The potential can be made stationary for $\sigma_1 = \sigma_2 = 0$ by requiring

$$F_1(\xi) = 0 \tag{14a}$$

from which follows

$$F_2(\xi) = 2\xi F_3(\xi) \tag{14b}$$

The cosmological constant can be avoided if the condition

$$F_0(\xi) = 0 \tag{14c}$$

is imposed. Then (13a) is put into the form

$$\mathcal{L}_1 = -b \tfrac{1}{2} \mu^2 \tilde{\sigma}_2{}^2 + O(\sigma^3) \tag{15a}$$

where the diagonalized fields are defined by

$$\tilde{\sigma}_1 = (1 + \xi^2 Z)^{-1/2} (\sigma_1 + \xi Z^{1/2} \sigma_2) \tag{15b}$$

$$\tilde{\sigma}_2 = (1 + \xi^2 Z)^{-1/2} (\xi Z^{1/2} \sigma_1 - \sigma_2) \tag{15c}$$

and the squared mass of $\tilde{\sigma}_2$ is given by

$$\mu^2 = 2 f^2 \kappa^{-2} (1 + \xi^2 Z) F_3(\xi) \tag{15d}$$

We assume that the right-hand side of (15d) is positive under the constraints (14) for ξ and c's. The field $\tilde{\sigma}_1$ remains massless.

The same procedure can also be applied to $\mathcal{L}_2$. We find

$$\mathcal{L}_2 \longrightarrow b\{-M\overline{\Psi}\Psi + [\xi(3h_2 + 2\xi^{-1}h_1) Z^{1/2}\sigma_1 - h_2\sigma_2]\, \overline{\Psi}\Psi - \tfrac{1}{2} m^2 \Phi^2$$
$$+ [v_2 \xi (2g_2^2 + \xi^{-2} g_1^2) Z^{1/2} \sigma_1 - g_2^2 v_2 \sigma_2]\, \Phi^2 + O(\sigma^2)\} \tag{16a}$$

where

$$M = v_2 (h_2 + \xi^{-1} h_1) \tag{16b}$$

$$m^2 = v_2{}^2 (g_2{}^2 + \xi^{-2} g_1^2) \tag{16c}$$

give the masses of Ψ and Φ, respectively.

We assume that M or m is of the order of GeV, a typical mass scale of elementary particles; we must allow a latitude of a factor $\sim 10^5$ to accommodate a wide spectrum ranging from the electron to the intermediate weak bosons. We have

$$\kappa M \sim 10^{-19} \tag{17a}$$

We also recall that, as was shown in Reference 3, results which are reasonable from an overall point of view emerge if we choose

$$f \sim 1, \quad \xi \sim 1, \quad \kappa v_1 \sim \kappa v_2 \sim 1$$
$$g_1 \sim g_2 \sim h_1 \sim h_2 \sim 10^{-19} \tag{17b}$$

The second and the fourth terms in (16a) hence describe "gravitational" interactions of the matter fields characterized by a small constant κ. This justifies ignoring the higher-order terms in (16a).

The Yukawa interaction term is now given the form

$$\mathcal{L}_3 \longrightarrow -be(1 - 3\gamma\sigma_1)\,\overline{\Psi}\Psi\Phi + O(\sigma^2) \tag{18}$$

§(4): *Sources of the Scalar Fields*

It is convenient to introduce an effective Lagrangian

$$L = -\overline{\Psi}(\partial\!\!\!/ + M)\Psi - \tfrac{1}{2}\Phi_{,\mu}\Phi^{,\mu} - \tfrac{1}{2}m^2\Phi^2 - e\overline{\Psi}\Psi\Phi \tag{19}$$

for the matter fields which have acquired the masses as given by (16b) and (16c), with the other gravitational effects being neglected. From (19) follow the effective field equations,

$$(\partial\!\!\!/ + M)\Psi + e\Psi\Phi = 0 \tag{20a}$$

$$-(\Box - m^2)\Phi + e\overline{\Psi}\Psi = 0 \tag{20b}$$

From (19) we can also define the effective (symmetrized) matter energy-momentum tensor $T_{\mu\nu}$. By using (20), the trace T is calculated to be[2]

$$T = -M\overline{\Psi}\Psi - m^2\Phi^2 - \tfrac{1}{2}\Box\Phi^2 \tag{21}$$

The interaction term does not occur in (21), basically because T measures violation of scale invariance.

In order to calculate the sources of $\tilde{\sigma}$'s to the first order in κ, we use (20) in the sum of $\mathcal{L}_2$, $\mathcal{L}_3$, and in the second and third terms in the square bracket of (12), keeping terms up to linear in σ's. We also use the formula

$$\Phi_{,\mu}\Phi^{,\mu} = \tfrac{1}{2}\Box\Phi^2 - \Phi\Box\Phi \tag{22}$$

[2] We can modify the theory slightly so that the last term of (21) is absent (the improved energy-momentum tensor), but without any substantial changes in the final results.

and discard the first term $\Box\Phi^2$, which is unimportant in the final results.[3] We obtain

$$\mathcal{L}' \approx b[\tfrac{1}{2} e\overline{\Psi}\Psi\Phi + (1 + \xi^2 Z)^{1/2} (h_2 \overline{\Psi}\Psi + g_2^2 v_2 \Phi^2)\, \tilde{\sigma}_2] \tag{23}$$

Notice the absence of $\tilde{\sigma}_1$. The sources $\tilde{J}_i$ of $\tilde{\sigma}_i$ are obtained:

$$\tilde{J}_1 = 0 \tag{24a}$$

$$\tilde{J}_2 = -(1 + \xi^2 Z)^{1/2} (h_2 \overline{\Psi}\Psi + g_2{}^2 v_2 \Phi^2) \tag{24b}$$

As explained in Section I, composition independence for rigid and static macroscopic objects will be true if $\tilde{J}_2$ is proportional to the first two terms of (21). (The last term is again neglected for the same reason as in $\tilde{J}_2$.) The proportionality follows if

$$M/h_2 = m^2 / v_2 g_2^2 \tag{25a}$$

or, on using (16b) and (16c),

$$h_1/h_2 = \xi^{-1} (g_1^2/g_2^2) \tag{25b}$$

It is straightforward to extend the analysis to a system of many Dirac fields and scalar fields in mutual interactions through Yukawa couplings of arbitrary combinations. We can show that (24a) continues to hold while (24b) is replaced by

$$\tilde{J}_2 = -(1 + \xi^2 Z)^{1/2} \left(\sum_\alpha h_{2(\alpha)} \overline{\Psi}_{(\alpha)} \Psi_{(\alpha)} + v_2 \sum_\beta g_{2(\beta)}^{\,2} \Phi_{(\beta)}^{\,2} \right) \tag{26}$$

where the subscripts in the parentheses label the kinds of fields. The right-hand side of (21) is also replaced by the corresponding sum. The condition (25b) for composition independence is now put into the form

$$\frac{h_{1(1)}}{h_{2(1)}} = \frac{h_{1(2)}}{h_{2(2)}} = \cdots = \xi^{-1} \frac{g_{1(1)}^{\,2}}{g_{2(1)}^{\,2}} = \xi^{-1} \frac{g_{1(2)}^{\,2}}{g_{2(2)}^{\,2}} = \cdots \tag{27}$$

The obviously simplest and plausible choice is

$$h_{1(\alpha)} = g_{1(\beta)} = 0 \tag{28}$$

which implies "specialized roles" of the scalar fields; apart from the mutual interaction $\mathcal{L}_1$, ϕ couples exclusively to the scalar curvature whereas ψ couples primarily to the matter fields.

[3] The term $\frac{1}{2}\Box\Phi^2$ contributes the additional terms,

$$\tilde{J}_1' = -\tfrac{1}{2}\kappa f^{-1} Z^{1/2} (1 + \xi^2 Z)^{-1/2} \Box\Phi^2$$

$$\tilde{J}_2' = -\tfrac{1}{2}\kappa^2 f^{-2} Z v_2 (1 + \xi^2 Z)^{-1/2} \Box\Phi^2$$

In the weak-field limit, $\Box\tilde{\sigma}_1 = \tilde{J}_1' \sim \Box\Phi^2$ gives $\tilde{\sigma}_1$ which is "frozen" inside the matter distribution Φ^2. In $\tilde{J}_2'$ we write $\Box\Phi^2 = (\Box - \mu^2)\Phi^2 + \mu^2\Phi^2$. The first part again yields the frozen component, while the second part is $\sim\mu^2/m^2 \sim 10^{-39}$ as small as $\tilde{J}_2$.

For the choice (28) the proportionality is given the form (apart from the terms containing $\Box\Phi_{(\beta)}^2$)

$$\tilde{J}_2 = v_2^{-1}(1+\xi^2 Z)^{1/2}\ T \tag{29}$$

The static gravitational potential is simply the sum of the usual long-range part and the finite-range part due to the exchange of $\tilde{\sigma}_2$. From (29) we find

$$a = G_\infty^{-1}\left[(1+\xi^2 Z)/4\pi v_2^2\right] = 2\xi^{-2} f^{-2}(1+\xi^2 Z) \tag{30}$$

in agreement with equation (13) of Reference 3 with $\zeta = 1, x = \xi$.

§(5): *Discussion*

The analyses in the preceding sections may be generalized further to any matter systems described by the theories with dimensionless coupling constants and the masses generated by v's. Various types of gauge theories, e.g., Weinberg–Salam model, will fall into this category, although the relation between Higgs fields in these theories and the scalar fields, especially ψ, in the present theory is yet to be worked out. It seems likely that only the Higgs fields acquire masses through the mechanism discussed here[4]; masses of other fields may emerge as a secondary effect from the vacuum expectation values of the Higgs fields.

One may be suspicious if the effect of interactions, like atomic or nuclear binding energies, is not included correctly in (21), which contains no interaction terms explicitly. We point out, however, that the terms in (21) do not represent energies either. Deviations of (21) from the energies of free constituent particles at rest should give exactly the interaction energies under the condition that the spatial integrals of $T_{\mu\nu}$ other than T_{00} vanish. Notice also that the fields are interacting fields obeying (20).

Our theory is somewhat unique in that we started with the (globally) scale-invariant total Lagrangian. In some other versions of the scalar–tensor theory, the particle masses are either introduced at the outset [8, 9] or generated spontaneously but with some dimensional constants already prepared in the potential term [6]. Perhaps the real origin of our nonvanishing v's will be found in "subtraction points" which are required to *define* quantum loop corrections in completely massless theories.[5] This is another aspect that distinguishes our approach from the purely classical theory in which dimensional constants are introduced in the Lagrangian as coefficients of terms allowed from general requirements [9].

Although these conjectures on the origin of the dimensional constants do

[4] A negative squared mass can be obtained if the sign of the second term of (2e) is reversed.
[5] A similar view is also shared by Minkowski [5], in his scale-invariant theory.

not affect most of phenomenological consequences, our suggestion [2, 3] $\lambda^{-2} \sim Gm_N^4 \hbar^{-3} c$ does contain Planck's constant.[6]

Finally we add a comment on the formal argument on the equivalence principle. The finite-range part $V_2 = -aG_\infty Me^{-\mu r}/r$ coming from the $\tilde{\sigma}_2$ exchange has no origin in the space-time geometry. It thus appears that a test particle falls off a geodesic. We find, however, that $V_2(r)$ is equivalent mathematically to the contribution of a "position-dependent mass," as discussed by Brans and Dicke [8], and is absorbed into the "geometrical" term $V_1 = -\frac{1}{2}(1 + g_{00})$ by redefining the metric as $g_{00} \to g_{00} - 2V_2$. This can be achieved by a conformal transformation $g_{\mu\nu} \to (1 + 2V_2) g_{\mu\nu}$ [8]. A test particle then falls along a geodesic in the new geometry.

References

1. Long, D. R. (1976). *Nature*, **260**, 417; Stacey, F. D., Tuck, G. J., Maher, A. R., and Morris, D. (1978). *Geophys. Res. Lett.*, **5**, 377; Yu, H. T., Ni, W. T., Hu, C. C., Liu, F. H., Yang, C. H., and Liu, W. N. (1979). *Phys. Rev. D*, **20**, 1813; Hirakawa, H., Tsubono, K., and Oide, K. (1980). *Nature*, **283**, 184; Spero, R., Hoskins, J. K., Newman, R., Pellam, J., and Shultz, J. (1980). *Phys. Rev. Lett.*, **44**, 1645; Panov, V. I., and Frontov, V. M. (1979). *Sov. Phys. JETP*, **50**, 852; Paik, Ho Jung. (1979). *Phys. Rev. D*, **19**, 2320; Mills, A. P., Jr. (1979). *Gen. Rel. Grav.*, **11**, 1. See also Sugimoto, D. (1972). *Prog. Theoret. Phys.*, **48**, 699; Fujimoto, M., and Sugimoto, D. (1972). *Prog. Theoret. Phys.*, **48**, 705; Mikkelson, D. R., and Newman, M. J. (1977). *Phys. Rev. D*, **16**, 919; Yamaguchi, Y. (1977). *Prog. Theoret. Phys.*, **58**, 723; Blinnikov, S. I. (1978). *Astrophys. Space Sci.*, **59**, 13; Gibbons, G. W., and Whiting, B. F. (1980). DAMPT preprint.
2. Fujii, Y. (1971). *Nature Phys. Sci.*, **234**, 5; O'Hanlon, J. (1972). *Phys. Rev. Lett.*, **29**, 137; Acharya, R., and Hogan, P. A. (1973). *Nuovo Cimento Lett.*, **6**, 668.
3. Fujii, Y. (1974). *Phys. Rev. D*, **9**, 874.
4. Fujii, Y. (1972). *Ann. Phys. (N.Y.)*, **69**, 494; (1975). *Gen. Rel. Grav.*, **6**, 29.
5. Minkowski, P. (1977). *Phys. Lett.*, **71B**, 419.
6. Zee, A. (1979). *Phys. Rev. Lett.*, **42**, 417.
7. Roll, P. G., Krotkov, K., Dicke, R. H. (1964). *Ann. Phys. (N.Y.)*, **26**, 442; Braginsky, V. B., and Panov, V. I. (1971). *Sov. Phys. JETP*, **34**, 464.
8. Brans, C., and Dicke, R. H. (1961). *Phys. Rev.*, **124**, 925.
9. Wagoner, R. V. (1970). *Phys. Rev. D*, **1**, 3209.
10. Deser, S., and Zumino, B. (1977). *Phys. Rev. Lett.*, **38**, 1433.
11. Fujii, Y., and Nishino, H. (1979). *Z. Phys. C*, **2**, 247.
12. Scherk, J. (1979). *Phys. Lett.*, **88B**, 265.

[6] The same distance was also suggested by S. Deser and B. Zumino [10] (see also Y. Fujii and H. Nishino [11]) and by J. Scherk [12]. It seems natural to expect that some of the lower-spin fields in extended supergravity theories may result in the gravitational effects similar to the one discussed here.