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FROM THE WEYL THEORY TO A THEORY OF LOCALLY ANISOTROPIC SPACE-TIME

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It is shown that Weyl's ideas, pertaining to local conformal invariance, find natural embodiment within the framework of a relativistic theory based on a viable Finslerian model of space-time. This is associated with the peculiar property of the Finslerian metric which describes a locally anisotropic space of events. Such a metric, in contrast to the Riemannian one, is a conformal invariant, in which case the local conformal transformations of the Riemannian metric tensor, apart from space-time intervals, leave invariant rest masses as well as all observables and thus appear as local gauge transformations. The corresponding Finslerian theory of gravitation turns out, as a result, to be an Abelian gauge theory. It satisfies the principle of correspondence with Einstein's theory and predicts a number of nontrivial physical effects accessible for experimental test under laboratory conditions.

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Unified field theories, based on the supersymmetric generalization of the Weyl theory of gravitation [1], are distinguished by the fact that infinities are eliminated from diagrams with any number of loops. This remarkable feature of these theories, as well as difficulties involved in them, are largely inherited from the original conformally invariant Weyl theory. It will be shown in this paper that the idea of local conformal invariance can be raised to a level of local gauge invariance and thus, in place of the Weyl theory, one can construct a physically reasonable classical Abelian gauge theory within the framework of which the local conformal transformations of the Riemannian metric tensor appear as truly gauge transformations, i.e. transformations which leave invariant all the observables including rest masses and space-time intervals. The corresponding viable theory, preserving the merits of the Weyl theory, turns out to be devoid of its disadvantages and, so, appears to be very promising from the standpoint of quantization.

Before turning to the formulation of a new theory it is expedient to direct our attention to the pioneer studies of Weyl [2, 3]. Trying to construct a unified theory of gravitation and electromagnetism, Weyl proceeded from the requirement of invariance of the theory with respect to the local conformal transformations of the Riemannian metric tensor

$$g_{ik} \rightarrow e^{2\sigma(x)} g_{ik} \quad . \tag{1}$$

The appearing Abelian gauge field, denoted below by A_i, was identified by Weyl with an electromagnetic field. The basis for such identification was, evidently, the gradient character of the corresponding gauge transformations

$$A_i \rightarrow A_i + \alpha\, \sigma_{;i} \quad , \tag{2}$$

where a is a constant with a length dimension. It is easy, however, to show that in reality in the Weyl theory the A_i is not an electromagnetic field but acquires mass due to the violation of local conformal invariance. This result is essentially contained in the Weyl variational principle

$$\delta\left[\frac{1}{16\Lambda}\int\left(R+\frac{6}{a}A^{i}_{;i}-\frac{6}{a^2}A_iA^i\right)^2\sqrt{-g}\,d^4x-\frac{1}{4}\int F_{ik}F^{ik}\sqrt{-g}\,d^4x\right]=0\,, \qquad (3)$$

where $F_{ik}=A_{k;i}-A_{i;k}$ and Λ is a constant.

In fact, using the Weyl gauge condition

$$R+\frac{6}{a}A^{i}_{;i}-\frac{6}{a^2}A_iA^i=-4\Lambda \qquad (4)$$

and Eq. (3), we arrive at the following system of equations

$$F^{ik}{}_{;k}-\frac{6}{a^2}A^i=0\quad , \qquad (5)$$

$$R_{ik}-\frac{1}{2}Rg_{ik}-\Lambda g_{ik}=-F_{il}F^{l}{}_{k}+\frac{1}{4}F_{lm}F^{lm}g_{ik}+\frac{6}{a^2}\left(A_iA_k-\frac{1}{2}A_mA^mg_{ik}\right)\,. \qquad (6)$$

Since in virtue of (5) the equality $A^{i}_{;i}=0$ holds and in virtue of (6) the equality $R-\frac{6}{a^2}A_iA^i=-4\Lambda$ holds, it is clear that the gauge condition (4) is fulfilled for the solutions of the system (5) and (6). Thus the system of the fourth-order field equations, following from the gauge invariant variational principle (3), together with Eq. (4) fixing the gauge, is equivalent to the system of equations (5) and (6). In other words, the fixing of the gauge, which implies in the Weyl theory the incorporation of a length and mass standard, converts the A_i into a massive vector field and reduces the equations for gravitational field to the Einstein equations with the Λ term.

The above observation once again indicates that the general relativity theory describes the world in which local conformal symmetry has been violated. And the conventional way, which turned out to be not very fruitful, of reestablishing such symmetry consists in simultaneous local conformal transformations of space-time intervals and all rest masses [4]. In this case the corresponding transformations by the very meaning are not local gauge transformations since they affect the observables. This disadvantage of the theory can be eliminated by the only method - by assuming that the true geometry of space of events differs from the Riemann one and a new metric is so dependent on g_{ik} that the conformal transformations (1) leave invariant space-time intervals and are thus indeed local gauge transformations. Although at first sight this idea seems to be too daring, there are additional arguments in favour of it which, in combination with Weyl's ideas, make it possible to construct an alternative relativistic theory of space-time and gravitation.

The assumption that the geometry of space-time differs from the Riemann one is equivalent to the assumption that a group of local relativistic symmetry differs from a Lorentz group. This assumption is not void of sense since it is not unlikely [5] that the Lorentz transformations, experimentally verified for relatively small Lorentz factors, no longer work at relative velocities extremely close to the velocity of light. For the first time a grave suspicion about violation of the Lorentz transformations at the Lorentz factors $\gamma \gtrsim 5 \times 10^{10}$ appeared in connection the experimental study of the spectra of superhigh-energy primary cosmic protons. If this suspicion is confirmed in subsequent experiments [6] then this will signify that in nature other, generalized Lorentz transformations are realized which correctly relate the coordinates of events at any values of the Lorentz factor. The generalized Lorentz transformations should, evidently, possess the following properties: (i) be linear, (ii) constitute a group, (iii) provide invariance of the electrodynamic equations, (iv) lead to Einstein's law of the addition of three-velocities, and (v) be markedly different from the usual Lorentz transformations only at velocities extremely close to the velocity of light. It appeared [7] that the corresponding transformations do exist and their invariant is a metric

$$ds^2 = \left[\frac{(dx_0 - \vec{\nu}d\vec{x})^2}{dx_0^2 - d\vec{x}^2}\right]^{\tau} (dx_0^2 - d\vec{x}^2) \quad . \tag{7}$$

Representing a concrete homogeneous function for the differentials of coordinates of the second degree of homogeneity, this metric belongs to the type of Finslerian metrics [8]. It generalizes the Minkowski metric and describes a flat anisotropic space of events. Anisotropy is characterized by two constant invariant parameters: a parameter τ , fixing its value, and a unit vector $\vec{\nu}$ indicating the preferred direction in three-space. In this case to the isotropic Minkowski space there corresponds a value $\tau = 0$.

Perhaps the most pronounced physical manifestation of space anisotropy consists in that the daily rate of a moving clock with respect to mutually stationary synchronized clocks depends not only on the value but also on the direction of its velocity. At certain velocities the moving clock will even lead those at rest. Yet, on returning to the initial point, it will necessarily be slow. It is precisely owing to this circumstance that the notion of rectilinear inertial motion is conserved in a space with constant anisotropy and the action for a free particle turns out to be minimal on a straight world line.

No less impressive is the dependence of kinetic energy on the direction of particle velocity. Since the kinetic energy is equal to the work to overcome the inertial force, done when a particle is accelerated to a given velocity, it thus appears that in different directions a particle of mass m resists acceleration differently. Thus the effective inertial mass of a particle in an anisotropic space becomes a tensor explicitly dependent on the parameters r and $\vec{\nu}$ which determine the space anisotropy

$$m_{\alpha\beta} = m(1-r)(\delta_{\alpha\beta} + r\nu_\alpha\nu_\beta) . \quad (8)$$

If now this result is confronted with fundamental Mach's idea, according to which the inertial mass of a particle just should be a tensorial quantity dependent on the distribution and motion of external matter (and, hence, represent a tensor field on space-time), then such confrontation suggests that the parameters r and $\vec{\nu}$ should be considered not constants but fields on space-time, having distributed matter as their source. It means, in turn, that we should introduce into consideration a space-time with anisotropy varying from point to point. Then due to the dependence on the fields r and $\vec{\nu}$, which now determine a locally anisotropic space-time, the inertial mass will acquire the character of a tensor field corresponding to Mach's principle.

We thus arrive at a curved locally anisotropic space-time with the metric

$$ds = \left[\frac{(\nu_i dx^i)^2}{g_{ik}dx^i dx^k} \right]^{r/2} \sqrt{g_{ik}dx^i dx^k} \quad (9)$$

In contrast to the Riemann space, the Finsler space (9) has at different points different flat tangent spaces which differ by the values of r and $\vec{\nu}$. In this case the whole family of tangent spaces is described by Eq. (7).

Apart from the Riemannian metric tensor g_{ik} , related to gravitational field, the geometry of the Finsler space (9) depends on two additional fields: a scalar field τ , which determines the magni tude of local anisotropy, and a zero-vector field ν_i which indicate the locally preferred directions in space-time. A complex of the interacting fields g_{ik} , τ and ν_i forms, along with matter, a unified dynamic system, the quantitative description of which is involved in the task of the Finslerian theory of gravitation.

Owing to the structure of the metric (9), the anisotropic field directly affects only the motion of test bodies, which is described by the equations of Finslerian geodesics [9]. As regards the propagation of massless particles it occurs along the zero geodesics of a *Riemann space with the metric tensor* g_{ik} *. Therefore for the* synchronization of clocks and for reestablishing the three-geometry in a Finsler space of events the procedure of exchange of light signals is still applicable. As a result, taking into account the relationship between the local time $d\tau$ and the coordinate time dx^0 , i.e, the relation

$$d\tau = \frac{1}{c}\left(\frac{\nu_0^2}{g_{00}}\right)^{\tau/2}\sqrt{g_{00}}\,dx^0 , \qquad (10)$$

we obtain [10]

$$dl^2 = \left(\frac{\nu_0^2}{g_{00}}\right)^{\tau}\left(-g_{\alpha\beta} + \frac{g_{0\alpha}g_{0\beta}}{g_{00}}\right)dx^\alpha dx^\beta , \qquad (11)$$

$$\Delta x^0 = -g_{0\alpha}dx^\alpha / g_{00} . \qquad (12)$$

In the Riemannian case, i.e. where the field $\tau = 0$, the conformal transformations (1) change space-time intervals whereas for a Finsler space of events there exists invariance of the corresponding observables (11) and (10) under the simultaneous transformations

$$g_{ik} \to e^{2\sigma(x)} g_{ik} \quad , \quad \nu_i \to e^{\frac{(\tau-1)\sigma(x)}{\tau}} \nu_i \quad . \tag{13}$$

The transformations (13) also leave invariant the Finslerian metric (9) and the equations of motion of massive particles. In this case the rest masses remain constant. Thus within the framework of the Finslerian theory the local transformations (13) turn out to be truly gauge transformations.

The requirement of local gauge invariance plays the key role in generalizing the formalism of the theory of gravitation with the inclusion of anisotropy of the space of events. The dynamical equations of the generalized theory and all the observables should be constaructed so as to be independent of the concrete gauge of the fields g_{ik} and ν_i . For example, the gauge invariant action for compressible liquid in a locally anisotropic space appears as

$$S = -\frac{1}{c}\int \overset{*}{\mu} \left(\frac{\nu_i v^i}{\sqrt{g_{ik} v^i v^k}} \right)^{4\tau} \sqrt{-g}\; d^4x \quad , \tag{14}$$

where $\overset{*}{\mu}$ is the invariant liquid energy density, $v^i = dx^i/ds$, and ds is the Finslerian metric (9).

In order to describe the self-consistent dynamics of the Finsler space (9) and of compressible liquid, it is also necessary to formulate the gauge invariant variational principle for the fields g_{ik} , ν_i and τ . For this purpose one should first extend the given dynamical system by including in it (in accordance with (13)) a pair of Abelian gauge fields A_i and B_i. Then, however, it can be somewhat narrowed down, prescribing the Weyl field A_i as purely gauge, i.e. as a gradient of a scalar composed of g_{ik} , ν_i and τ . This makes it possible to confine our consideration to a reduced dynamical system with the only gauge field B_i. Its gradient transformations have the form

$$B_i \longrightarrow B_i + b\left[\frac{(\tau-1)\sigma(x)}{\tau}\right]_{;i} \qquad (15)$$

where b is a constant with a length dimension. As a result, the dynamics of the reduced system are described by the following variational principle

$$\delta\int\Big\{-\frac{1}{2}[\cdots]R-\frac{3}{4}[\cdots]^{-1}[\cdots]^{;i}[\cdots]_{;i}-\frac{\tau^{;i}\tau_{;i}}{4\varsigma(\varepsilon-\tau)}\left(\frac{\cdots}{\cdots}\right)^{2\tau}-\frac{f}{4}\mathcal{N}_{ik}\mathcal{N}^{ik}\left(\frac{\cdots}{\cdots}\right)^{2\tau-2}+$$
$$+\frac{\lambda^2 f}{2}\vartheta_i\vartheta^i\left(\frac{\cdots}{\cdots}\right)^{4\tau-2}-\frac{1}{4}\mathcal{F}_{ik}\mathcal{F}^{ik}-\frac{8\pi\hat{k}}{c^4}\overset{*}{\mu}\left(\frac{\vartheta_iV^i}{\sqrt{V_kV^k}}\right)^{4\tau}\Big\}\sqrt{-g}\,d^4x=0\,, \qquad (16)$$

where $\left(\frac{\cdots}{\cdots}\right)=\left(\vartheta^k\tau_{;k}/\sqrt{-\tau_{;k}\tau^{;k}}\right)$, $[\cdots]=\left[(1-\tau/\varepsilon)\left(\frac{\cdots}{\cdots}\right)^{2\tau}\right]$. $\mathcal{N}_{ik}=\vartheta_{k;i}-\vartheta_{i;k}-(\vartheta_kB_i-\vartheta_iB_k)/b$,

$\mathcal{F}_{ik}=B_{k;i}-B_{i;k}$, R is a Riemannian scalar, the constants f, ς and $1/\varepsilon$ are dimensionless, in which case the ς determines the interaction of matter with the field τ and the $1/\varepsilon$ determines the interaction of the field τ with g_{ik}, $\hat{k}$ is a starting gravitational constant related to the observed Newtonian constant k by the relation $\hat{k}=k/\eta$, η is a renormalization constant given by the expression

$$\eta = 1+\frac{\varsigma/(2\varepsilon)}{\sqrt{1+\varsigma/(4\varepsilon)}}$$

and, finally, λ^2 is a Lagrangian multiplier.

The variational principle (16) leads to the hydrodynamic equations [11] and to a gauge invariant system of field equations. In the gauge, fixed by the condition $\vartheta^k\tau_{;k}=\sqrt{-\tau^{;k}\tau_{;k}}$, the system of field equations is of the form

$$-\left(1-\frac{\tau}{\varepsilon}\right)\left(R_{ik}-\frac{1}{2}Rg_{ik}\right)-\frac{1}{\varepsilon}\left(\tau_{;i;k}-\tau^{;n}_{;n}g_{ik}\right)-\frac{\varepsilon+3\varsigma}{2\varsigma\varepsilon(\varepsilon-\tau)}\left(\tau_{;i}\tau_{;k}-\frac{1}{2}\tau_{;n}\tau^{;n}g_{ik}\right)-$$
$$-\mathcal{F}_{il}\mathcal{F}_k{}^{l}+\frac{1}{4}\mathcal{F}_{lm}\mathcal{F}^{lm}g_{ik}+f\left(-\mathcal{N}_{il}\mathcal{N}_k{}^{l}+\frac{1}{4}\mathcal{N}_{lm}\mathcal{N}^{lm}g_{ik}\right)+f\lambda^2\vartheta_i\vartheta_k+$$
$$+f\left[\frac{1}{2}\mathcal{N}_{lm}\mathcal{N}^{lm}+\frac{8\pi k}{c^4 f\eta}\tau\left(\overset{*}{\mu}-3p\right)\left(\frac{\vartheta_lV^l}{\sqrt{V_nV^n}}\right)^{4\tau}\right]\left[\frac{\vartheta_i\tau_{;k}+\vartheta_k\tau_{;i}}{\vartheta^n\tau_{;n}}-\frac{\tau_{;i}\tau_{;k}}{\tau^{;n}\tau_{;n}}\right]+$$
$$+\frac{8\pi k}{c^4\eta}\left\{\left[(1-\tau)\overset{*}{\mu}+(1+3\tau)p\right]\frac{V_iV_k}{V_nV^n}-pg_{ik}\right\}\left(\frac{\vartheta_lV^l}{\sqrt{V_nV^n}}\right)^{4\tau}=0\,, \qquad (17)$$

$$-\tau^{;\kappa}_{;\kappa}+\frac{\varsigma f}{2}\mathcal{N}_{\ell m}\mathcal{N}^{\ell m}-$$
$$-2\varsigma f(\varepsilon-\tau)\Big\{\Big[\frac{1}{2}\mathcal{N}_{\ell m}\mathcal{N}^{\ell m}+\frac{8\pi k}{c^4 f\eta}\tau(\overset{*}{\mu}-3p)\Big(\frac{\nu_\ell V^\ell}{\sqrt{V_n V^n}}\Big)^{4\tau}\Big]\Big(\frac{\tau^{;\kappa}}{\tau_{;n}\tau^{;n}}-\frac{\nu^\kappa}{\tau_{;n}\nu^n}\Big)\Big\}_{;\kappa}+$$
$$+\frac{8\pi k\varsigma}{c^4\eta}(\overset{*}{\mu}-3p)\Big(\frac{\nu_\ell V^\ell}{\sqrt{V_n V^n}}\Big)^{4\tau}\Big[1+(\varepsilon-\tau)\ln\frac{(\nu_\ell V^\ell)^2}{V_n V^n}\Big]=0\ , \tag{18}$$

$$\mathcal{N}^{i\kappa}{}_{;\kappa}+\frac{1}{8}\mathcal{N}^{i\kappa}B_\kappa-\lambda^2\nu^i-$$
$$-\Big[\frac{1}{2}\mathcal{N}_{\ell m}\mathcal{N}^{\ell m}+\frac{8\pi k}{c^4 f\eta}\tau(\overset{*}{\mu}-3p)\Big(\frac{\nu_\ell V^\ell}{\sqrt{V_n V^n}}\Big)^{4\tau}\Big]\frac{\tau^{;i}}{\sqrt{-\tau_{;n}\tau^{;n}}}+$$
$$+\frac{8\pi k}{c^4 f\eta}\tau(\overset{*}{\mu}-3p)\Big(\frac{\nu_\ell V^\ell}{\sqrt{V_n V^n}}\Big)^{4\tau}\frac{V^i}{\nu_\ell V^\ell}=0\ , \tag{19}$$

$$\mathcal{F}^{i\kappa}{}_{;\kappa}-\frac{f}{8}\mathcal{N}^{i\kappa}\nu_\kappa=0\ , \tag{20}$$

$$\nu_\kappa\nu^\kappa=0\ , \tag{21}$$

where p is pressure.

In order to estimate the interaction constants ς and $1/\varepsilon$, occurring in these equations, it is necessary to proceed as follows. First, on obtaining the static centrally symmetric solution of the field equations (17-21), i.e. on solving the Finslerian Schwarzschild problem, it is necessary to integrate the equations of Finslerian geodesics. Whereupon, on finding corrections to the classical gravitation effects due to the local anisotropy of space, produced by the Sun, it is necessary to require that these corrections be not inconsistent with experimental data relating to the solar system. As a result [12], we get

$$- 0.054 \lesssim \varsigma \leqslant 0; \; 0 < 1/\varepsilon \lesssim 0.25 \qquad (22)$$

Within the framework of the theory outlined here the equality $\varsigma = 0$ indicates the absence of the field τ which determines the magnitude of the local anisotropy of space-time. In this case the Finslerian metric (9) reduces to the Riemannian one and the Finslerian theory of gravitation, to Einstein's theory. And if $\varsigma \neq 0$ then the existence of local anisotropy in space-time leads to nontrivial physical consequences. For example, in virtue of Eq. (8), the dynamical properties of a particle, present in anisotropic space, turn out to be similar to those of a quasiparticle in a crystal. In both cases the effective mass is a tensorial quantity and, according to Newton's second law, the acceleration direction does not coincide, generally speaking, with the direction of an applied force. Analogy between the local anisotropy of space and the crystalline state of a medium is enhanced by the fact that the appearance of local anisotropy is energetically favourable and may be regarded as a peculiar phase transition in the geometrical structure of space-time, viz. the transition from the "amorphous" Riemannian to the "ordered" Finslerian structure. However, since the major source of a local anisotropy field is a trace of the energy-momentum tensor for the fields of matter, the violation of the local isotropy of space-time does not occur, apparently, spontaneously but turns out to be a secondary effect caused by the spontaneous violation of gauge symmetry in some or other version of unified theory when fields acquire masses and the energy-momentum tensor acquires a nonzero trace. As a result, the high-temperature phase transitions with the violation of gauge symmetries are accompanied by the corresponding phase transitions in the geometric structure of space-time. In this case the geometric transitions determine the most energetically favourable channels of violation of gauge symmetries. This circumstance makes very promising the application of the theory of locally anisotropic space-time in modern cosmological models.

The above-mentioned limitation on the interaction constants makes it possible to conclude that the upper bound of the space local anisotropy, produced by the Sun near the Earth's orbit, is at a level of 10^{-10}. If the anisotropy were not less than 5×10^{-10} then it could be revealed in the laboratory experiments [13] carried out as long as 20 years ago with a view to verifying the relativity theory. Up to date, in connection with the development of engineering, it has become possible to enhance the resolution of such experiments by three orders of magnitude and thereby to give them a status of experiments on the search for local anisotropy of space. The setting-up of corresponding search experiments becomes especially urgent at present in view of the fact that within the framework of the relativistic theory of locally anisotropic space-time we have a strong alternative to the postulate of rigorous local isotropy of space - the fundamental postulate of Einstein's relativity theory.

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Георгий Юрьевич БОГОСЛОВСКИЙ

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