Graviweak interactions and their role in gravitational dynamics and electrodynamics

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(Submitted 8 September 1994)
Zh. Eksp. Teor. Fiz. **107**, 283–298 (February 1995)

Lagrangian densities that include terms with broken P and C symmetry and that lift the degeneracy of spin states are postulated for spinor particles and photons in a gravitational field. We obtain modified Dirac and Maxwell equations, and construct semiclassical Hamiltonians for particles and macroscopic objects that possess intrinsic quantized and classical angular momentum, respectively. We show that graviweak interactions can lead to a number of effects, including altered object weight due to partial alignment of electron and nucleon spins, angular splitting and relative delay between left- and right-hand polarized waves passing near the sun, modification of atomic spectra and the structure of atomic shells, generation of energetic x-rays near neutron stars, and the advent of isotropically distributed high-energy cosmic rays at an early stage in the evolution of the universe. © *1995 American Institute of Physics.*

1. INTRODUCTION

The interactions of fermions and bosons with the gravitational field in current theories of gravitation are degenerate with respect to spin variables, and are invariant under charge conjugation, space inversion, and time reversal. There is no such degeneracy over spin states in purely electromagnetic interactions, although the interactions themselves remain invariant under C-, P-, and T-transformations. The weak interactions (now merged with the electromagnetic interactions) comprise another well-known class with broken C- and P- symmetry.

The present stock of experimental and theoretical data provides no serious basis for entertaining the notion that gravitational interactions, in contrast to the others, ought to remain degenerate over spin states and be invariant under C-, P-, and T-. These ideas led the present author to postulate[1] a generalized structure for the Lagrangian density $\mathscr{L}_1$ of spinor particles in a gravitational field that lifted the degeneracy over spin states and contained terms with broken C, P, and T symmetry. If, however, the new gravitational interactions were universal, i.e., if they applied equally well to fermions and bosons, then in addition to the modified density $\mathscr{L}_1$ there should exist a modified Lagrangian density $\mathscr{L}_2$ for, say, photons ($\mathscr{L}_2$ did not appear in Ref. 1). In either case, the introduction of C-, P-, and T-breaking terms that also lifted the spin degeneracy ought to be a single operation, in a certain sense (invoking the same constants, for example), that should not lead to inconsistencies with the semiclassical relationship between energy and momentum obtained under the same conditions using one density or the other.

We show below that the hypothetical graviweak interactions are consistent with the available experimental data. If we assume that these interactions are responsible for a number of observed effects that are as yet unexplained, then by comparing the theoretical results with the experimental, we can determine the coupling constant for the new interactions. Many of the effects induced by these interactions should then become more tangible, and deserving of the attention of experimentalists. Several experiments are entirely feasible under typical terrestrial conditions.

2. GENERAL ASSUMPTIONS

We postulate a simpler Lagrangian density $\mathscr{L}_1$ than in Ref. 1 for spinor particles in a gravitational field.[1)]

$$\mathscr{L}_1 \equiv \frac{i\hbar}{2}\sqrt{-g}\Big\{\bar\psi\Gamma^\alpha(\nabla_\alpha\psi)-(\nabla_\alpha\bar\psi)\Gamma^\alpha\psi+\frac{C}{8}\bar\psi\gamma^5[(\Gamma^\alpha\Gamma^\beta-\Gamma^\beta\Gamma^\alpha)(\nabla_\alpha\Gamma_\beta-\nabla_\beta\Gamma_\alpha)-(\nabla_\alpha\Gamma_\beta-\nabla_\beta\Gamma_\alpha)\times(\Gamma^\alpha\Gamma^\beta-\Gamma^\beta\Gamma^\alpha)]\psi\Big\}-m\sqrt{-g}\bar\psi\psi. \quad (1)$$

Here $\Gamma^\alpha \equiv h^\alpha_{(a)}\gamma^a$, the $h^\alpha_{(a)}$ are comoving tetrads ($h^\alpha_{(a)}=g^{\alpha\beta}h_{(a)\beta}$, $h^{(a)\alpha}=\eta^{ab}h^\alpha_{(b)}$, etc.), the $g_{\alpha\beta}$ are the metric coefficients in Riemannian space (with no curvature), $\eta^{ab}|_{a\neq b}=0$, $\eta^{ab}|_{a=b}=(1,-1,-1,-1)$, and the matrices γ^a with Roman indices ($a=0,1,2,3$) and Γ^α with Greek indices ($\alpha=0,1,2,3$) satisfy the relations

$$\gamma^a\gamma^b+\gamma^b\gamma^a=2\eta^{ab}, \quad \Gamma^\alpha\Gamma^\beta+\Gamma^\beta\Gamma^\alpha=2g^{\alpha\beta},$$

with

$$\nabla_\alpha\psi=\partial_\alpha\psi+\frac{1}{4}(\nabla_\alpha h^\beta_{(a)})h_{(b)\beta}\gamma^b\gamma^a\psi,$$

$$\nabla_\alpha h_{(a)\beta}=\partial_\alpha h_{(a)\beta}-\Gamma^\nu_{\alpha\beta}h_{(a)\nu},$$

where $\Gamma^\nu_{\alpha\beta}$ is a Christoffel symbol. The graviweak coupling constant C is real by virtue of the hermiticity of $\mathscr{L}_1$. Clearly, the density $\mathscr{L}_1$ takes the familiar form when $C=0$, and in the absence of gravitational fields, it reduces to the Lagrangian density for free spinor particles.

Let the source of the gravitational field be a static, spherically symmetric object of mass $M \gg m$. Then to leading order in M—to which we adhere throughout—we have in Galilean coordinates

$$g_{00}=1-\frac{2M}{r},\quad g_{kn}\Big|_{k=n}=-1-\frac{2M}{r},\quad g_{kn}\Big|_{k\neq n}=g_{0n},$$
$$=0,$$

$$h^0_{(0)}=1+\frac{M}{r},\quad h^k_{(n)}\Big|_{k=n}=1-\frac{M}{r},\quad h^k_{(0)}=h^k_{(n)}\Big|_{k\neq n}=0.$$

Inserting these expressions into (1), we obtain

$$(\gamma^a\hat{P}_a-m)\psi=0, \tag{2}$$

in which

$$\hat{P}_0\equiv\left(1+\frac{M}{r}\right)E,\quad \hat{\mathbf{P}}\equiv\frac{1}{2}\left\{\left(1-\frac{M}{r}\right)\hat{\mathbf{p}}+\hat{\mathbf{p}}\left(1+-\frac{M}{r}\right)\right\}$$
$$-i\gamma^5C\frac{M\hbar\mathbf{r}}{2r^3},$$

E determines the energy of the spinor particle, and the normalization equation for ψ is

$$\int\sqrt{-g}d^3x\bar{\psi}\Gamma^0\psi=1. \tag{3}$$

For subsequent use, we can also write Eq. (2) in the form

$$P_0\psi=(\boldsymbol{\alpha}\hat{\mathbf{P}}+m\rho_3)\psi\equiv\hat{H}_0\psi. \tag{2a}$$

With the normalization (3), it can easily be shown that to leading order in M, the expectation value of any observable $\hat{Q}$ in the state ψ and its time derivative are given by

$$\langle\hat{Q}\rangle=\frac{1}{2}\int\sqrt{-g}d^3x\psi^+\left[\left(1+\frac{M}{r}\right),\hat{Q}\right]_+\psi,$$
$$\frac{d\langle\hat{Q}\rangle}{dt}=\left\langle\frac{\partial\hat{Q}}{\partial t}\right\rangle+\frac{i}{\hbar}\int\sqrt{-g}d^3x\psi^+\left\{[\hat{H}_0,\hat{Q}]_-+\frac{1}{2}\left[\hat{H}_0,\left[\hat{Q},\frac{M}{r}\right]_-\right]_+\right\}\psi, \tag{4}$$

where $[...,...]_-$ and $[...,...]_+$ denote a commutator and an anticommutator.

Noting that

$$P_0\boldsymbol{\sigma}=\frac{1}{2}(\hat{H}_0\boldsymbol{\sigma}+\boldsymbol{\sigma}\hat{H}_0)=m\rho_3\boldsymbol{\sigma}+\rho_1\hat{\mathbf{P}},$$
$$P_0\rho_1=\frac{1}{2}(\hat{H}_0\rho_1+\rho_1\hat{H}_0)=(\boldsymbol{\sigma}\hat{\mathbf{P}}), \tag{5}$$

we introduce the covariant definition of the four-spin in the semiclassical approximation,

$$\mathbf{s}=\frac{\hbar}{2m}\{m\langle\rho_3\boldsymbol{\sigma}\rangle+\langle\rho_1\rangle\mathbf{P}\},\quad s_0=\frac{(\mathbf{sP})}{P_0},$$

where $\mathbf{P}\equiv[1-(M/r)]\langle\hat{\mathbf{p}}\rangle$, and $\langle...\rangle$ denotes an average over a localized wave packet. In this approximation, the semiclassical quantities $\mathbf{s}$, s_0, $\mathbf{P}$, and P_0 will by virtue of (4) vary according to

$$\frac{d\mathbf{s}}{dt}\equiv\dot{\mathbf{s}}=-\frac{Ms_0\mathbf{r}}{r^3}-\frac{M}{P_0r^3}[\mathbf{sL}],\quad \dot{s}_0=-\frac{M(\mathbf{sr})}{r^3},$$
$$\dot{\mathbf{P}}=-\frac{MP_0\mathbf{r}}{r^3}-\frac{M}{P_0r^3}[\mathbf{PL}],\quad \dot{P}_0=-\frac{M(\mathbf{rP})}{r^3}, \tag{6}$$

where $\mathbf{L}$ is the angular momentum of an object with classical spin. The resulting relations

$$P_0^2-\mathbf{P}^2=m^2,\quad s_0^2-\mathbf{s}^2=\text{const}\equiv-\boldsymbol{\zeta}^2$$

enable us to express s_0 and $\mathbf{s}$ in terms of the three-vector $\boldsymbol{\zeta}$:

$$\mathbf{s}=\boldsymbol{\zeta}+\frac{s_0\mathbf{P}}{P_0+m},\quad s_0=\frac{(\boldsymbol{\zeta}\mathbf{P})}{m}.$$

Hence (6),

$$\dot{\boldsymbol{\zeta}}=-\frac{2P_0+m}{P_0(P_0+m)}\frac{M[\boldsymbol{\zeta}\mathbf{L}]}{r^3},$$

which then yields the expressions for the spin precession of a nonrelativistic and an ultrarelativistic top:

$$\dot{\boldsymbol{\zeta}}=-\frac{3M}{2mr^3}[\boldsymbol{\zeta}\mathbf{L}],\quad \dot{\boldsymbol{\zeta}}=-\frac{2M}{Er^3}[\boldsymbol{\zeta}\mathbf{L}]. \tag{7}$$

As expected, by virtue of their quantum nature, graviweak interactions do not appear in the classical expressions for the precession of tops in a gravitational field. Their influence on the dynamics of a classical spin system shows up—albeit extremely weakly—when quantum effects are taken into consideration [see (33) and (34)]. In elementary particle applications, Eqs. (6) and (7) cannot hold in general, since the very notion of a "classical" spin is inadmissible. To find the equations of motion of the spin $\mathbf{s}$ of a particle with nonzero rest mass, we revert to (4) and (5), retaining quantum terms in $\hat{H}_0$ and $\hat{\mathbf{P}}$, yielding

$$\dot{\mathbf{s}}=-\frac{Ms_0\mathbf{r}}{r^3}-\frac{M}{P_0r^3}[\mathbf{sL}]+\frac{CM}{r^3}[\mathbf{sr}]+\frac{CMs_0}{P_0r^3}\mathbf{L},$$
$$\dot{s}_0=-\frac{M(\mathbf{sr})}{r^3}+\frac{CM}{P_0r^3}(\mathbf{sL}), \tag{6a}$$

whereupon

$$\dot{\boldsymbol{\zeta}}=-\frac{2P_0+m}{P_0(P_0+m)}\frac{M[\boldsymbol{\zeta}\mathbf{L}]}{r^3}+\frac{CMm[\boldsymbol{\zeta}\mathbf{r}]}{P_0r^3}+\frac{CM(\mathbf{rP})[\boldsymbol{\zeta}\mathbf{P}]}{P_0(P_0+m)r^3}. \tag{7a}$$

We see that the spin of a particle with nonzero rest mass precesses under the influence of graviweak interactions not only about $\mathbf{L}$, but about the particle's radius vector $\mathbf{r}$ and velocity vector $\mathbf{P}/P_0$.

If the particle is massless, then by (5), $\mathbf{s}=\zeta(\hbar/2)\times(\mathbf{P}/P_0)$, where $\zeta=\pm1$, i.e.,

$$\dot{\mathbf{s}}=-\frac{2M}{Er^3}[\mathbf{sL}]. \tag{7b}$$

Limiting attention to leading order in $\hbar$ and M, and to positive-energy states (where the rest mass is included), we can go from (2) to the equations for two-component spinors χ. Taking the usual approach, we obtain

$$E\chi=\left\{\left(1-\frac{3M}{r}\right)\frac{\hat{\mathbf{p}}^2}{2m}-\frac{mM}{r}+\frac{3M\hbar}{4mr^3}(\boldsymbol{\sigma}'\hat{\mathbf{L}})-\frac{CM\hbar}{2r^3}(\boldsymbol{\sigma}'\mathbf{r})\right\}\chi \tag{8}$$

for nonrelativistic particles, where E is the particle's total energy minus its rest mass, and

$$\left(1+\frac{2M}{r}\right)E^2\chi=\left\{\left(1-\frac{2M}{r}\right)\hat{\mathbf{p}}^2+\frac{2M\hbar}{r^3}(\boldsymbol{\sigma}'\hat{\mathbf{L}})-\frac{CM\hbar}{Er^3}(\mathbf{r}\hat{\mathbf{p}})[\boldsymbol{\sigma}'\hat{\mathbf{p}})\right\}\chi \tag{9}$$

for a particle with $m=0$. In Eqs. (8) and (9), $\hat{\mathbf{L}}\equiv[\mathbf{r}\hat{\mathbf{p}}]$, $\boldsymbol{\sigma}'$ represents the Pauli spin matrices, and the normalization condition for the two-component spinors is

$$\int\sqrt{-g}d^3x\chi^+\chi=1.$$

Taking the semiclassical limit in (8) by averaging over a localized wave packet, we introduce the Hamiltonian[2)]

$$H=\left(1-\frac{3M}{r}\right)\frac{\mathbf{p}^2}{2m}-\frac{mM}{r}+\frac{3M}{2mr^3}(\mathbf{sL})-\frac{CM}{r^3}(\mathbf{s}_2\mathbf{r}) \tag{10}$$

for a nonrelativistic object with classical and total quantum spins ($\mathbf{s}_1$) and ($\mathbf{s}_2$) ($\mathbf{s}=\mathbf{s}_1+\mathbf{s}_2$). Here the total quantum spin $\mathbf{s}_2$ will be proportional to N_s, the number of spin-aligned fermions in the object: $\mathbf{s}_2=(\hbar/2)N_s\mathbf{s}_2^0$, where $\mathbf{s}_2^0$ specifies the spin alignment direction.

Similarly going to the semiclassical limit in (9), we obtain the Hamiltonian

$$H=\left(1-\frac{2M}{r}\right)\left\{\mathbf{p}^2+\frac{4M}{r^3}(\mathbf{sL})-\frac{2CM}{pr^3}(\mathbf{rp})(\mathbf{sp})\right\}^{1/2}. \tag{11}$$

Noting that to leading order in $\hbar$ and M the spin of a massless particle remains unaltered when it is collinear with the particle's velocity, i.e., $\mathbf{s}=\zeta(\hbar/2)\mathbf{n}$, where $\zeta=\pm1$ and $\mathbf{n}\simeq\mathbf{p}/p$, one can easily derive the spin precession equation (7b) from (11), as well as the energy–momentum relation

$$E=\left(1-\frac{2M}{r}\right)\left\{\mathbf{p}^2-\zeta C\frac{2Ms}{r^3}(\mathbf{rp})\right\}^{1/2}. \tag{12}$$

It would be fair to assume that graviweak interactions ought to be just as universal as conventional gravitational interactions—i.e., they ought to apply to bosons as well as fermions. If we then take $s=\hbar$ in (12), it should apply equally well to photons. This necessitates a modification of the electromagnetic field Lagrangian $\mathscr{L}_2$:

$$\mathscr{L}_2\equiv\frac{\sqrt{-g}}{8\pi}(g^{\alpha\mu}g^{\beta\nu}-g^{\alpha\nu}g^{\beta\mu})\times[(A^{(+)}_{\alpha,\beta}+A^{(-)}_{\alpha,\beta})-2i\zeta Ch^{(a)}_{\alpha,\beta}h^{\sigma}_{(a)}(A^{(+)}_{\sigma}-A^{(-)}_{\sigma})]^+\times[(A^{(+)}_{\mu,\nu}+A^{(-)}_{\mu,\nu})-2i\zeta Ch^{(b)}_{\mu,\nu}h^{\tau}_{(b)}(A^{(+)}_{\tau}-A^{(-)}_{\tau})]. \tag{13}$$

Here $A^{(+)}_\alpha$ and $A^{(-)}_\alpha$ are the positive- and negative-frequency parts of the potential $A_\alpha\equiv A^{(+)}_\alpha+A^{(-)}_\alpha$;

$$\zeta=s_{[kn]0}T_{m0}e^{0knm}|s_{[pq]0}T_{l0}e^{0pql}|^{-1}=\pm1$$

if the numerator is nonzero, and $\zeta=0$ otherwise. For electromagnetic waves, $\zeta=0$ corresponds to linear polarization and $\zeta=\pm1$ to right- and left-handed circular polarization (rcp, lcp).

From (13), we then have

$$g^{\alpha\beta}A^{\lambda}_{;\alpha;\beta}-g^{\beta\lambda}A^{\alpha}_{;\beta;\alpha}-2i\zeta C[(g^{\lambda\mu}g^{\alpha\nu}-g^{\lambda\nu}g^{\alpha\mu})h^{(a)}_{\mu;\nu}h^{\tau}_{(a)}(A^{(+)}_{\tau}-A^{(-)}_{\tau})]_{;\alpha}=\frac{4\pi}{\sqrt{-g}}j^{\lambda}, \tag{14}$$

where semicolons denote covariant derivatives, and the classical current density is

$$j^{\lambda}=\sum_i e_i\frac{dx^{\lambda}}{dx^0}\delta(\mathbf{r}-\mathbf{r}_i).$$

For rcp ($\zeta=1$) and lcp ($\zeta=-1$) waves, the source-free solution of (14) is, to first order in M and $\hbar$,

$$\mathbf{A}=\frac{1}{\sqrt{2}}\sum_{\boldsymbol{\kappa}}A(\boldsymbol{\kappa})\left\{(\mathbf{a}_1+i\zeta\mathbf{a}_2)\exp\left(-i\omega t+i\int\boldsymbol{\kappa}d\mathbf{r}\right)+(\mathbf{a}_1-i\zeta\mathbf{a}_2)\exp\left(i\omega t-i\int\boldsymbol{\kappa}d\mathbf{r}\right)\right\}, \tag{15}$$

where $A(\boldsymbol{\kappa})$ is the amplitude of a monochromatic wave, $\mathbf{a}_1(\boldsymbol{\kappa})$ and $\mathbf{a}_2(\boldsymbol{\kappa})$ are mutually orthogonal unit vectors that specify the wave polarization (these are orthogonal to $\boldsymbol{\kappa}$ as well), and $\boldsymbol{\kappa}$ satisfies the dispersion relation

$$\left(1+\frac{2M}{r}\right)\omega^2-\left(1-\frac{2M}{r}\right)\boldsymbol{\kappa}^2+\frac{2\zeta CM}{r^3}(\mathbf{r}\boldsymbol{\kappa})=0. \tag{16}$$

Clearly (16) is the same as (12). Thus, the new graviweak interaction can be introduced in a consistent manner.

Finally, let us consider a fermion of mass m that carries charge e. The external fields are electromagnetic (with four-potential $A^{\alpha}\equiv\Phi,\mathbf{A}$) and gravitational, the latter being produced both by a static, stationary centrosymmetric object of mass $M\gg m$ and by other sources, including the electromagnetic field itself. We neglect gravitational fields of all sources except M. Incorporating the interaction with the electromagnetic field $A^{\alpha}\equiv\Phi,\mathbf{A}$ into the Lagrangian density (1), we obtain to leading order in M

$$(\gamma^a \hat{P}_a - m)\psi = 0,$$

$$\hat{P}_0 \equiv \left(1+\frac{M}{r}\right) i\hbar\partial_0 - \left(1-\frac{M}{r}\right) e\Phi, \qquad (17)$$

$$\hat{\mathbf{P}} \equiv \frac{1}{2}\left\{\left(1-\frac{M}{r}\right)\hat{\mathbf{p}} + \hat{\mathbf{p}}\left(1-\frac{M}{r}\right)\right\} - \left(1+\frac{M}{r}\right) e\mathbf{A} - i\gamma^5 \frac{CM\hbar\mathbf{r}}{2r^3}.$$

This equation can be compared with the equation for a one-electron atom in an external field with $\Phi = \Phi_1 + \Phi_2$, where Φ_1 is the scalar potential that the nucleus at $\mathbf{r}_0$ produces at the electron at $\mathbf{r}$, and Φ_2 and $\mathbf{A}$ are the electromagnetic field potentials outside the atom. Granted, Eq. (17) then describes a system in which the atomic nucleus is stationary, and if the atom is moving relative to the external fields, (17) will have the correct form only insofar as the nuclear state can be considered constant in the region of interest. Small variations in the nuclear state due to external fields can in principle be addressed by perturbation methods, as is done, for example, in allowing for nuclear recoil due to interactions with the electron.

Shifting now to the quasirelativistic approximation in (17), which takes quartic terms in the velocity into account, we obtain the following equation to first order in $\hbar M$ and neglecting terms of order $(\cdots)a_0 M/r_0^2$, where a_0 is the linear size of the atom:

$$\left\{\left(1+\frac{3M}{r_0}\right) i\hbar\partial_0 + \left(1+\frac{2M}{r_0}\right)\frac{mM}{r_0}\right\}\chi$$
$$= \left\{\frac{\hat{\mathbf{p}}^2}{2m} + \left(1+\frac{M}{r_0}\right)e\Phi - \left(1+\frac{2M}{r_0}\right)\frac{e}{2m}[(\mathbf{A}\hat{\mathbf{p}})+(\hat{\mathbf{p}}\mathbf{A})]\right.$$
$$+ \left(1+\frac{4M}{r_0}\right)\frac{e^2\mathbf{A}^2}{2m} - \left(1+\frac{2M}{r_0}\right)(\hat{\boldsymbol{\mu}}[\nabla\mathbf{A}])$$
$$+ \frac{e\hbar}{4m^2}(\boldsymbol{\sigma}'[(\partial_0\mathbf{A}+\nabla\Phi)\hat{\mathbf{P}}_1]) + \frac{e\hbar^2}{8m^2}\nabla^2\Phi$$
$$- \frac{(\boldsymbol{\sigma}'\hat{\mathbf{P}}_1)^4}{8m^3} + \frac{M\hbar e}{m^2 r_0}(\boldsymbol{\sigma}'[\partial_0\mathbf{A}\hat{\mathbf{P}}_1])$$
$$+ \frac{3M\hbar}{4mr_0^2}(\boldsymbol{\sigma}'[\mathbf{n}_3\hat{\mathbf{p}}]) + \frac{M\hbar e}{4mr_0^2}(\boldsymbol{\sigma}'[\mathbf{n}_3\mathbf{A}])$$
$$\left. + \frac{M\hbar e\Phi}{4m^2 r_0^2}(\boldsymbol{\sigma}'[\mathbf{n}_3\hat{\mathbf{P}}_1]) - \frac{CM\hbar}{2r_0^2}\sigma_3'\right\}\chi. \qquad (18)$$

Here

$$\hat{\boldsymbol{\mu}} \equiv \frac{e\hbar}{2m}\boldsymbol{\sigma}', \quad \hat{\mathbf{P}}_1 \equiv \left(1-\frac{M}{r_0}\right)\hat{\mathbf{p}} - \left(1+\frac{M}{r_0}\right)e\mathbf{A},$$

and the 3 axis is taken to be parallel to $\mathbf{r}_0$ and $\mathbf{n}_3 \equiv \mathbf{r}_0/r_0$.

Clearly, the postulated graviweak interactions lead to a modification of all fundamental equations governing both fermions and bosons, and this comes about in a consistent manner. We next consider the consequences of this new form of gravitational interaction, and possible detection experiments.

3. DELAY AND ANGULAR SPLITTING EFFECTS IN LEFT- AND RIGHT-HAND CIRCULARLY POLARIZED BEAMS PROPAGATING IN THE FIELD OF A STATIC, CENTROSYMMETRIC GRAVITATIONAL SOURCE

By virtue of the spin-state dependence of the Hamiltonian (11), we expect differences in the propagation of left- and right-hand polarized photons in the gravitational field of a source M. We now consider the dynamics in more detail.

Using canonical variables $\mathbf{r}$ and $\mathbf{p}$, we obtain from (11)

$$\mathbf{v} = \frac{1}{E}\left(1-\frac{4M}{r}\right)\left\{\mathbf{p} + \frac{2M}{r^3}[\mathbf{sr}] - \zeta\frac{CM\hbar\mathbf{r}}{r^3}\right\}. \qquad (19)$$

Reexamining (11) with this in mind, and working once again to leading order in $\hbar$ and M, we have

$$\left(1+\frac{4M}{r}\right)\mathbf{v}^2 = 1, \qquad (20)$$

which is equivalent to requiring that the interval between two events on a world line vanish: $ds^2 \equiv g_{\alpha\beta}dx^\alpha dx^\beta = 0$. Further bearing in mind that according to (19)

$$\mathbf{L} = [\mathbf{rp}] = \left(1+\frac{4M}{r}\right)E[\mathbf{rv}] - \frac{2M}{r^3}[\mathbf{r}[\mathbf{sr}]],$$

and therefore

$$E^2 r^4(\dot{\theta}^2 + \dot{\varphi}^2\sin^2\theta) = \left(1-\frac{4M}{r}\right)\left[1-\frac{4M}{r}(1-q)\right]\mathbf{L}^2,$$

where $q \equiv (\mathbf{sL})/\mathbf{L}^2$, we have from (20) that[3)]

$$\dot{r}^2 = \left(1-\frac{4M}{r}\right)\left\{1-\frac{\mathbf{L}^2}{E^2 r^2}\left(1-\frac{4M}{r}\right)\right\}. \qquad (21)$$

Both E and $\mathbf{L}^2$ are constant; the latter comes from

$$\dot{\mathbf{L}} = -\dot{\mathbf{s}} = \frac{2M}{Er^3}[\mathbf{sL}].$$

According to (21), the pericenter of the photon trajectory at $\varphi = 0, r = r_0$ is given by

$$r_0^2 = \frac{\mathbf{L}^2}{E^2}\left(1-\frac{4M}{r_0}\right). \qquad (22)$$

Plugging (22) into (21), we then have

$$\dot{r} = \pm\sqrt{\frac{r^2-r_0^2}{r}}\left\{1-\frac{2M}{r}-\frac{2Mr_0}{(r_0+r)r}\right\}. \qquad (23)$$

It must be emphasized that r_0 is not the same for lcp and rcp photons. In fact, in a state with given initial momentum $\mathbf{p}_1$ and squared angular momentum $\mathbf{L}^2$, the energy E of a photon emitted at $r = r_1, \varphi = \varphi_1$ will be

$$E = p_1\left(1-\frac{2M}{r_1}\right)\left\{1+\zeta\frac{CM\lambda\sqrt{r_1^2-r_0^2}}{2\pi r_1^3}\right\}, \qquad (24)$$

where $\lambda \equiv 2\pi\hbar/p_1$ is the wavelength corresponding to the photon energy. Thus, according to (22),

$$r_0=\bar{r}_0\left\{1-\zeta C\,\frac{M\lambda\sqrt{r_1^2-\bar{r}_0^2}}{2\pi r_1^3}\right\}. \qquad (25)$$

If the source and the photon detector are located at points (r_1,φ_1) and (r_2,φ_2) on opposite sides of the trajectory's pericenter, then we find from (23) that the gravitational delay time Δt of photons with polarization ζ is

$$\Delta t=2M\,\ln[(r_2+\sqrt{r_2^2-r_0^2})(r_1+\sqrt{r_1^2-r_0^2})/r_0^2]. \qquad (26)$$

Invoking (25), we obtain the ratio of the gravitationally induced differential delay $\delta t\equiv\Delta t|_{\zeta=1}-\Delta t|_{\zeta=-1}$ between rcp and lcp photons and the mean gravitational delay $\bar{\Delta} t$,

$$\frac{\delta t}{\bar{\Delta} t}=\frac{2CM\lambda\sqrt{r_1^2-\bar{r}_0^2}}{\pi r_1^3}\,\ln^{-1}\frac{r_2+\sqrt{r_2^2-\bar{r}_0^2}}{r_1-\sqrt{r_1^2-\bar{r}_0^2}}. \qquad (27)$$

There is no such differential delay if the entire analysis is instead based on linear polarization states.

Before assessing feasibility, let us consider the characteristics of the trajectories of circularly polarized photons. To do so, and bearing in mind that deviations from the initial plane specified by the momentum $\mathbf{p}_1$ are negligible,[4)] we plug the relationship

$$\dot{\varphi}^2=\frac{r_0^2}{r^4}\left(1+\frac{4M}{r_0}-\frac{8M}{r}\right)$$

into (23). Integrating the latter, we obtain

$$\varphi=\arccos\frac{r_0}{r}+\frac{2M}{r_0}\sqrt{\frac{r-r_0}{r+r_0}}.$$

Hence, with $r_1,r_2\gg r_0$, we have for $\Delta\varphi$ angular deviation of photons with polarization ζ

$$\Delta\varphi=4M/r_0. \qquad (28)$$

The ratio of the differential deviation $\delta\varphi\equiv\Delta\varphi\times|_{\zeta=1}-\Delta\varphi|_{\zeta=-1}$ and the mean deviation $\overline{\Delta\varphi}$ will accordingly be

$$\frac{\delta\varphi}{\overline{\Delta\varphi}}=\frac{CM\lambda\sqrt{r_1^2-\bar{r}_0^2}}{\pi r_1^3}. \qquad (29)$$

There will thus be no angular splitting of radiation based on linear polarization states.

The data from the Pioneer 6 experiments that studied the propagation of polarized radio waves near the sun are well known.[5)] These suggested, with some degree of certainty, that the deviations of rcp and lcp rays were different in the solar gravitational field. For example, at a total deviation $\overline{\Delta\varphi}\sim0.2''$, the difference $\delta\varphi$ was approximately $0.002''$. At the same time, there was no noticeable differential deviation between linearly polarized waves.

If we assume that the observed effect was due to graviweak interactions, we can estimate the relevant coupling constant C. At the time of the observations, the distance r_1 from the sun to Pioneer 6 was of the order of the distance from Mars to the sun, and the observing wavelength[6)] was $\lambda\sim13$ cm. Then (29) yields $C\sim8\cdot10^{18}$ cm. If the gravitational delay is measured at $\bar{r}_0\sim4R_\odot$, then according to (27), the differential gravitational delay between rcp and lcp waves will be $\sim2\cdot10^{-3}$ times the total time delay $\overline{\Delta t}$ (measured at the same wavelength, $\lambda\sim13$ cm). This is consistent with the differential delay between rcp and lcp wave noted in the Pioneer 6 experimental results (see footnote 5). The differential delay between linearly polarized waves was essentially zero.

To improve our confidence in the consistency of an explanation based on graviweak interactions with the observed effects and to determine the coupling constant C more accurately, it would be desirable to repeat these experiments at a variety of wavelengths and at various spacecraft distances r_1 from the sun.

4. INFLUENCE OF GRAVIWEAK INTERACTIONS ON THE WEIGHT OF PARTIALLY SPIN-ALIGNED OBJECTS, AND THEIR NEGLIGIBLE ROLE IN SOLAR SYSTEM DYNAMICS

We see from (10) that the graviweak contribution makes interactions between an object and a gravitational source dependent upon the magnitude and direction of the object's aligned spin $\mathbf{s}_2=(\hbar/2)N_s\mathbf{s}_2^0$. Indeed, from (10) we have

$$\mathbf{v}=\left(1-\frac{3M}{r}\right)\frac{\mathbf{p}}{m}+\frac{3M}{2mr^3}[\mathbf{sr}], \qquad (30)$$

which then yields

$$\left(1+\frac{3M}{r}\right)m\dot{\mathbf{v}}=-\frac{mM\mathbf{r}}{r^3}\left[1+\frac{3}{2}\mathbf{v}^2-\frac{9(\mathbf{sL})}{2m^2r^2}+\frac{2C(\mathbf{s}_2\mathbf{r})}{mr^2}\right]+\frac{CM}{r^5}[\mathbf{r}[\mathbf{s}_2\mathbf{r}]]+\frac{3M}{r^3}\left\{[\mathbf{sv}]-\frac{3\dot{r}}{2r}[\mathbf{sr}]+\frac{1}{2}[\dot{\mathbf{s}}\mathbf{r}]+mr\dot{r}\mathbf{v}\right\}.$$

Hence, the weight of an object at rest becomes

$$F=-\left(1-\frac{3M}{r}\right)\frac{mM}{r^2}\left\{1+\frac{N_s}{N}\frac{C\hbar(\mathbf{s}_2^0\mathbf{r})}{m_\mathrm{p}r^2}\right\}, \qquad (31)$$

where we have made the replacement $m\simeq Nm_\mathrm{p}$ inside the curly brackets, so that N denotes the total number of nucleons in the object (recall that N_s is the number of spin-aligned fermions in the object).

If $N_s/N\sim0.1$ and the weighing takes place on earth, then the relative change in the weight of the object due to the second term in curly brackets, for $\mathbf{s}_2^0\mathbf{r}=\pm r$, will be $\sim3\cdot10^{-5}$. Obviously, if graviweak interactions actually exist, then under certain circumstances they can induce an appreciable change in an object's weight.

The author is aware of two papers[2,3] that deal with the weight (during damped rotation) of right-handed (in the terminology of Hayasaka and Takeuchi,[2] corresponding to $\mathbf{s}_1\mathbf{r}=-s_1r$) and left-handed ($\mathbf{s}_1\mathbf{r}=s_1r$) tops. Hayasaka and Takeuchi[2] claim that the right-handed top weighs less than a nonrotating top, while the left-handed top remains the same, with the relative change being a linear function of the top's angular velocity Ω ($|\delta F/F|\simeq4.6\cdot10^{-5}$ at $\Omega\sim10^3$ sec^{-1}).

No current theory of gravitation contains a mechanism capable of inducing such a huge relative change in weight, and an asymmetric one at that. Admittedly, no such phenomenon was detected by Fallar *et al.*,[3] i.e., the weights of the tops were always the same to within the errors, irrespective of the direction of rotation, even though the weighing precision was somewhat higher than that of Hayasaka and Takeuchi.[2] This of course casts the results in Ref. 2 into some doubt. Nevertheless, a close examination of the differences between the two experiments reveals the following. The rotors being weighed in Ref. 2 were made of metal (bronze, aluminum, steel), while those in Ref. 3 were nylon; those in Ref. 2 were initially spun up to $(14-15)\cdot 10^3$ rpm, and measurements began at $(12-13)\cdot 10^3$ rpm (with no further manipulation of the rotor after it was spun up), while in Ref. 3 it was initially spun up to $8\cdot 10^3$ rpm, and after some mechanical operations (hose removal, cap placement), measurements began at $6\cdot 10^3$ rpm; the rotor in Ref. 2 was in a vacuum and was spun up electromagnetically, while in Ref. 3 it was spun up by a gas jet through an attached hose, which can produce an overpressure condition in the chamber and heat the rotor.

If we suppose that the results of Hayasaka and Takeuchi[2] are in fact "clean," and that they are inherently due to graviweak interactions, then they can be accounted for by assuming that for some reason, under the conditions reported in Ref. 2, the spins of some fraction of the rotor's fermions were aligned during right-handed rotation in the same direction as the earth's rotation (i.e., parallel to the latter's angular momentum), and that that number gradually declined as the rotation slowed down. On the other hand, during left-handed rotation either too few spins were aligned, or none at all. Under the conditions of Ref. 3, alignment was either insufficient or completely lacking in both senses of rotation. If this is a valid interpretation, then (31) tells us that $N_s \simeq 0.22N$ at $12\cdot 10^3$ rpm, and $N_s \simeq 0.16N$ at $\Omega \sim 10^3$ sec^{-1}.

To assess just how "clean" the experiment of Ref. 2 may have been, and more to the point, the reality of graviweak interactions, it would be useful to mount the following experiment. Having prepared a cylindrical sample of some suitable material, a fair number of fermion spins would be aligned along its axis, and the sample would be weighed while standing on one end or the other. This would all be done, of course, at low temperatures in a thermally isolated nonmagnetic chamber, avoiding outside contamination when the cylinder was turned over, etc. If such an experiment were to yield negative results (at some high confidence level), the implication would be that the constant C is considerably smaller than the value obtained in Sec. 3, and that the phenomena discussed above (including those in Sec. 3) have some other explanation unrelated to graviweak interactions. If on the other hand the results were positive, they would directly confirm the existence of a new form of gravitational interaction in nature, which of course would be an exceptionally valuable outcome.

We now consider the effects of graviweak interactions on planetary dynamics in the solar system, which become significant due to the large value of the constant C. To do so, we make use of (10) and (30), which yield

$$E=\left(1+\frac{3M}{r}\right)\frac{m\dot{r}^2}{2}+\left(1-\frac{3M}{r}\right)\frac{\mathbf{L}^2}{2mr^2}-\frac{mM}{r}+\frac{3M}{2mr^3}(\mathbf{sL})-\frac{CM}{r^3}(\mathbf{s}_2\mathbf{r}) \tag{32}$$

and

$$\dot{\mathbf{L}}=\frac{3M}{2mr^3}[\mathbf{sL}]-\frac{CM}{r^3}[\mathbf{s}_2\mathbf{r}],$$
$$\dot{\mathbf{s}}_1=-\frac{3M}{2mr^3}[\mathbf{s}_1\mathbf{L}],\quad \dot{\mathbf{s}}_2=-\frac{3M}{2mr^3}[\mathbf{s}_2\mathbf{r}]+\frac{CM}{r^3}[\mathbf{s}_2\mathbf{r}]. \tag{33}$$

From the latter, we obtain in the present approximation the conservation law

$$\mathbf{L}^2-2CMm(\mathbf{s}_2\mathbf{n})=\text{const}\equiv\tilde{L}_0^2, \tag{34}$$

where $\mathbf{n}\equiv\mathbf{r}/r$. Substituting this into (32) yields

$$E=\left(1+\frac{3M}{r}\right)\frac{m\dot{r}^2}{2}-\frac{mM}{r}+\left(1-\frac{3M}{r}\right)\frac{\tilde{L}_0^2}{2mr^2}+\frac{3M}{2mr^3}(\mathbf{sL}). \tag{35}$$

We see from (33) that if the quantized aligned spin $\mathbf{s}_2$ is nonzero, it will precess slowly about $\mathbf{L}$ and very rapidly ($\omega_2 \simeq 2\cdot 10^8$ sec^{-1}) about $\mathbf{r}$. Therefore, only the component of $\mathbf{s}_2$ in the direction of $\mathbf{r}$ will effectively contribute to planetary dynamics. Even if we make the entirely improbable assumption that all of the fermion spins in the planet are either parallel or antiparallel to $\mathbf{r}$, the term containing $\mathbf{s}_2$ in (34) will be eight orders of magnitude less than $\mathbf{L}^2$. In actuality, N_s/N in planets is minuscule, so graviweak interactions have an exceedingly small role to play in planetary dynamics, and they can simply be neglected, putting $\tilde{L}_0^2 \simeq \mathbf{L}^2 \simeq \text{const}$, and $(\mathbf{sL}) \simeq (\mathbf{s}_1\mathbf{L})$.

5. INFLUENCE OF GRAVIWEAK INTERACTIONS ON ATOMIC SPECTRA IN A GRAVITATIONAL FIELD AND THE SPECTRUM OF ELECTROMAGNETIC RADIATION

Assuming the requisite conditions for (18) to hold, we make it our starting point, keeping in mind [see (14)] that the scalar potential $A^0\equiv\Phi$ of a motionless point nucleus with charge $-e$ will be given to first order in M by $\Phi(\mathbf{r}-\mathbf{r}_0)=-e/|\mathbf{r}-\mathbf{r}_0|$, where $\mathbf{r}_0$ is the radius vector of the nucleus and $\mathbf{r}$ is the observation point. Retaining only leading terms, we obtain in the motionless-nucleus approximation the following basic equation for atomic hydrogen:

$$\left(1+\frac{3M}{r_0}\right)E\chi=\left\{\frac{\hat{\mathbf{p}}^2}{2m}-\left(1+\frac{M}{r_0}\right)\frac{e^2}{r}-\frac{\hat{\mathbf{p}}^4}{8m^3}+\frac{\pi e^2\hbar^2}{2m^2}\delta(\mathbf{r})+\left(1-\frac{M}{r_0}\right)\frac{e^2\hbar}{4m^2r^3}(\boldsymbol{\sigma}'\hat{\mathbf{L}})-\frac{CM\hbar}{2r_0^2}\sigma_3'\right\}\chi, \tag{36}$$

in which $\mathbf{r}$ is the radius vector of the probable electron position relative to the nucleus, and $\hat{\mathbf{L}}\equiv[\mathbf{r}\hat{\mathbf{p}}]$.

The last term in (36) can (for example, in the vicinity of a neutron star) be comparable to or even greater than the Coulomb term (by several orders of magnitude). Therefore, it cannot be treated as a perturbation under those circumstances. Near the earth or sun, the term containing C is four to six orders of magnitude less than the Coulomb term, and therefore it can be treated as a perturbation.

We start by addressing the second situation, to which the zeroth-order solution is well known (see Ref. 4, for example), as are the correction to the energy spectrum due to the first three terms of the perturbation. The correction induced by the last term is, in the $j=l+\frac{1}{2}$ and $j=l-\frac{1}{2}$ states,

$$\Delta E_{n,j,m_j}^{(j=l+1/2)}=-\frac{CM\hbar m_j}{2jr_0^2},\quad \Delta E_{n,j,m_j}^{(j=l-1/2)}=\frac{CM\hbar m_j}{2(j+1)r_0^2}. \tag{37}$$

Consequently,

$$E_{n,j,m_j}^{(j=l+1/2)}=-\frac{\tilde{R}\hbar}{n^2}\left[1+\frac{\alpha^2}{n^2}\left(\frac{2n}{2j+1}-\frac{3}{4}\right)\right]-\frac{CM\hbar m_j}{2jr_0^2}, \tag{38}$$

where $\tilde{R}\equiv[1-(M/r_0)]me^4/2\hbar^3$.

Note that the factor $[1-(M/r_0)]$ in $\tilde{R}$ also leads to a gravitational frequency shift. To show this, let an observer located at $r_0=r_2$ receive photons emitted by atoms (in identical transitions) located at $r_0=r_2$ and $r_0=r_1$. By virtue of conservation of the photon energy $\hbar\omega$, a photon will arrive from $r_0=r_1$ at $r_0=r_2$ with energy $\hbar\omega_1\sim[1-(M/r_1)]$, while the energy of the photon emitted at $r_0=r_2$ will be proportional to $[1-(M/r_2)]$, so $\omega_1\neq\omega_2$. If $r_1<r_2$ then $\omega_1<\omega_2$, and $(\omega_1-\omega_2)/\omega_1\simeq-M(r_2-r_1)/r_1r_2$, i.e., the energy $\hbar\omega_1$ does not suffice to excite the atom. This then is none other than a redshift.

The corrections (37) lead to a splitting of the atomic energy levels according to the states m_j, and to the appearance of new spectral lines. One characteristic feature of this radiation is that it carries information about the mass M of the object near which it was produced, and about the distance r_0 of its source from the center of that object. Near the sun and the earth, radiation from the ground state of atomic hydrogen is produced at wavelengths $\lambda\simeq2.5\cdot10^{-2}$ cm and $\lambda\simeq0.74$ cm, respectively. To find the intensity of the radiation from the atomic ground states, we need to recognize that it can only be produced in spin-flip transitions. Since the operator σ_3' commutes with the perturbation $(\mathbf{A}_\gamma^+\hat{\mathbf{p}})$, where in accordance with (15) the second-quantized potential of the photons being produced is

$$\mathbf{A}_\gamma^+=\frac{1}{L^{3/2}}\sqrt{\frac{\pi\hbar}{\omega}}\hat{q}^+(\mathbf{a}_1-i\zeta\mathbf{a}_2)\exp\left(i\omega t-i\int\boldsymbol{\kappa}d\mathbf{r}\right),$$

this perturbation cannot produce spin-flip transitions—it cannot induce radiation.

The perturbation

$$\hat{V}=-\frac{e\hbar}{2m}(\boldsymbol{\sigma}'[\nabla\mathbf{A}_\gamma^+])$$

does not commute with σ_3', however, which leads to magnetic dipole radiation. Since graviweak interactions change the form of the relationship between $\boldsymbol{\kappa}$ and ω to that found in (16), the radiative transition probabilities then take the form

$$dw=\frac{2\pi}{\hbar}|\langle f|\hat{V}|i\rangle|^2\rho_f(\boldsymbol{\kappa}), \tag{39}$$

where

$$\rho_f=\left(\frac{L}{2\pi}\right)^3\frac{\kappa^2}{\hbar}\frac{d\kappa}{d\omega}d\Omega,$$

and with (37), to leading order in M,

$$\frac{d\kappa}{d\omega}=\frac{\kappa}{\omega}=\sqrt{1+\zeta^2\cos^2\theta}+\zeta\cos\theta,$$

where θ is the angle between $\boldsymbol{\kappa}$ and $\mathbf{r}_0$, and $\zeta=\pm1$. If instead we were discussing linear polarizations, we would have $d\kappa/d\omega=1$. The end result is that the angular distribution of the radiative intensity (to leading order in M) is

$$\frac{dI}{\sin\theta\,d\theta}=\frac{\alpha\hbar^3}{8m^2}\left(\frac{CM^4}{r_0^2}[\sqrt{1+\zeta^2\cos^2\theta}+\zeta\cos\theta]^5(1-\zeta\cos\theta)^2.\right. \tag{40}$$

Near the sun this yields an intensity of approximately $5\cdot10^{-21}$erg/sec.

We now turn to the case in which the term with C in (36) is no longer just a perturbation. Then the zeroth-order solution of Eq. (36) takes the form

$$\chi=\frac{1}{2}\begin{pmatrix}1+s\\1-s\end{pmatrix}R_{nl}(\rho)Y_l^m(\vartheta,\varphi), \tag{41}$$

where $R_{nl}(\rho)$ is a generalized Laguerre polynomial, $\rho\equiv2r/n\tilde{a}_0$, $\tilde{a}_0\equiv[1-(M/r_0)]a_0$, $a_0\equiv\hbar^2/me^2$, the $Y_l^m(\vartheta,\varphi)$ are spherical harmonics, and $s=\pm1$ corresponds to a positive or negative projection of the spin in the direction of $\mathbf{r}_0$. The atom's energy spectrum then becomes

$$E_{n,j,s}=-\frac{\tilde{R}\hbar}{n^2}\left[1+\frac{\alpha^2}{n^2}\left(\frac{2n}{2j+1}-\frac{3}{4}\right)\right]-\frac{CM\hbar s}{2r_0^2}. \tag{42}$$

Solutions of (41) and (42) exist for $E_{n,j,s}<0$, i.e., for $s=-1$ with $(M/r_0^2)<1.8\cdot10^{-13}(1/n^2)\text{cm}^{-1}$. If this condition fails to hold, then in multielectron atoms the spins will be preferentially aligned parallel to $\mathbf{r}_0$, and the old *aufbau* principle will be altered; accordingly, the spectrum will also change, due to the filling of higher levels. When $E_{n,j,-1}<0$, spin-flip transitions take place that give rise to radiation at frequencies

$$\omega_s=\frac{CM}{r_0^2}<4.3\cdot10^{16}\text{ s}^{-1}, \tag{43}$$

corresponding to $\lambda>4.4\cdot10^{-6}$ cm. As before, the intensity of this radiation will be governed by Eq. (40). At the highest admissible frequency, the radiated power will be of order $5\cdot10^{-6}$erg/sec.

Note that the expressions (43) and (40) characterize radiation both from nucleons in nuclei and from free fermions. In the latter case, if the fermion is relativistic, its intrinsic magnetic moment must be corrected in all calculations by replacing $\mu_0=e\hbar/2m$ with $\mu=\mu_0(m/E)$, which gives rise to a factor $(m/E)^2$ on the right-hand side of Eq. (40).

In neutron stars, $CM/r_0^2\lesssim3\cdot10^{11}\text{cm}^{-1}$. In the vicinity of such stars, it is therefore entirely possible that x rays or even gamma rays will be generated (with $\lambda\gtrsim2\cdot10^{-11}$). The radiative intensity can then reach values of order 10^{16}erg/cm if the fermions are nonrelativistic or weakly relativistic. The radiative intensity due to relativistic particles will be somewhat supressed by the relativistic factor.

Finally, there is one more interesting circumstance worth noting. We see from the structure of the Hamiltonian (10) that the term containing C is positive when $\mathbf{s}_2\mathbf{r}<0$, leading to repulsion of a fermion by the gravitational source M. If the magnitude of that repulsion exceeds the conventional gravitational attraction, such fermions will be accelerated away from M. All such particles (both e^- and p^+) will have left-handed polarization. Far from the source, they will have acquired an additional kinetic energy $\Delta E = CM\hbar/r_0^2$. Furthermore, this expression also follows directly from (2a), i.e., the form of ΔE is not a consequence of the nonrelativistic approximation. For typical neutron stars, $\Delta E \lesssim 10 m_e$.

The indicated mechanism for the acceleration of left-handed protons and electrons may provide one source of cosmic rays at that early stage in the evolution of the universe when fluctuations led to the formation of small-scale granular structure with granule masses M and linear dimensions a_0. If we assume that $1\,\mathrm{cm}^{-1} \lesssim M/a_0^2 \lesssim 10^6\,\mathrm{cm}^{-1}$, the energy ΔE acquired by left-handed fermions will lie in the range $(10^{14}-10^{20})$ eV. Larger granules (with smaller values of M/a_0^2) will produce left-handed fermions of lower energy. Due to the random spatial distribution of the granules and subsequent multiple scattering of particles, they should be distributed essentially isotropically throughout the universe. Incidentally, photons—even left-handed ones—are unaffected by this mechanism, since their coordinate velocity is always given by (20), which does not depend on the graviweak constant.

6. CONCLUSION

Standard theories of the interaction of matter (fermions and electromagnetic fields) with the gravitational field contain no reference to the spin states of the particles considered, i.e., they are degenerate in the spin variables. The theory proposed here lifts the degeneracy in that degree of freedom, and in a manner that subsumes both fermions and photons. The postulated graviweak interaction breaks both space and charge symmetry, endowing all of nature with such behavior, and not just electroweak processes.

In view of the serious consequences of the proposed theory, it would be highly desirable to mount a number of special experiments to test it; one such experiment (a direct test to determine C) has been outlined in Sec. 4. Bringing to bear the predictions of this theory, it would also be desirable to repeat investigations of the angular splitting of circularly polarized waves in the solar gravitational field.

The author thanks the Russian Fund for Fundamental Research for partial financial support of this work.

[1)] In units with $c=G=1$.
[2)] In the Hamiltonian (10) we have not included any term associated with the spin energy of the object, which is assumed to be small.
[3)] Since $\mathbf{s}\sim\mathbf{v}$, $\mathbf{sL}\simeq 0$ to the adopted approximation in M, and the term in q need no longer be retained.
[4)] The deviation of trajectories from the initial plane results from the conservation of the total momentum $\mathbf{J}=\mathbf{L}+\mathbf{s}$ and the curvature of the trajectories, i.e., from the fact that $\mathbf{s}$ rotates as the direction of $\mathbf{v}$ changes. It can easily be shown, however, that such deviations are exceedingly small, and we shall neglect them, taking $\dot{\theta}=0, \theta=\pi/2$.
[5)] Unfortunately, after the author became acquainted with this work, all of his relevant bibliographic data were lost (except for a few listings), and all attempts to recover them were unsuccessful. The citation of this source in the bibliography is therefore somewhat unconventional.
[6)] The radio wavelengths utilized by such spacecraft are typically ~3 cm and ~13 cm. The estimate of C requires refinement.

[1] Yu. M. Loskutov, Trudy Mat. Fiz. **91**, 309 (1992).
[2] H. Hayasaka and S. Takeuchi, Phys. Rev. Lett. **63**, 2701 (1989).
[3] J. E. Fallar, W. J. Hollander, P. G. Nelson, and M. P. McHugh, Phys. Rev. Lett. **64**, 825 (1990).
[4] A. A. Sokolov, Yu. M. Loskutov, and I. M. Ternov, *Quantum Mechanics*[in Russian], Prosveshchenie, Moscow (1965)[Holt, Rinehart, and Winston, New York (1966)].

Translated by Marc Damashek