WHAT DO WE KNOW IN FACT
ABOUT T-ODD, BUT P-EVEN INTERACTIONS?

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Abstract

General structure of fermion-fermion and photon-fermion interactions is presented. Radiative corrections, due to the P-odd part of the electroweak interaction, transform the T-odd, but P-even fermion-fermion interaction into a T-odd and P-odd one. The experimental information about T-odd, P-odd effects is sufficiently rich to obtain in this way new limits on the parameters of T-odd, P-even electron-electron, electron-nucleon and nucleon-nucleon interactions, as well as on some β-decay parameters. These limits are much better than those known previously.

INTRODUCTION

I met Art Rich only twice, at ICAP-11 in Paris in 1988, and just two years ago here in Ann Arbor. Certainly, other people can tell much more about this bright man and brilliant physicist. But let me also say here some words about him.

Our last meeting was about two weeks before his going to the hospital. In retrospect, when recalling some his jokes, it seems to me that Art was already aware of the serious operation ahead. And now I can appreciate his dignified behaviour that October.

It so happened also that only by chance, just before his death I could fully appreciate Art's warmth and kindness. It was too late already to reciprocate them.

I wish to think that Art would find the subject of this talk of some interest.

Direct experimental information on the T-odd, P-even (TOPE) interactions is rather poor. Best limits on the relative magnitude of the corresponding admixtures to nuclear forces lie around 10^{-3} [1-4]. We will relate below all interactions to the Fermi weak interaction constant G. Since the nuclear scale of weak interactions is $Gm_\pi^2 \sim 2 \times 10^{-7}$, those limits can be formulated as $10^4 G$. An experiment [5], going on now, aims at improving these limits by three orders of magnitude.

Experimental information on TOPE electron-nucleon interaction is practically absent. In ref. [6] an atomic experiment was suggested which can hopefully reach an accuracy about $\sim 3 \times 10^4 G$ (see also ref. [7]). Higher accuracy is aimed at in the recent experimental proposal [8].

As to TOPE electron-electron interaction, its possible manifestations in positronium were discussed previously in ref. [9].

Finally, upper limits on the relative magnitude of TOPE correlations in β-decay lie at best around 10^{-3} [10-15].

The main result of the present talk based on refs. [16-18] is that experimental data on the T- and P-odd effects lead to new, very strict upper limits on the TOPE electron-electron, electron-nucleon and nucleon-nucleon interactions, as well as on some β-decay parameters.

T-ODD, P-EVEN FERMION-FERMION AMPLITUDES

Let us start from the construction of the TOPE scattering amplitude for fermions of spin 1/2. To find the number of these amplitudes it is convenient to go over to the annihilation channel and to classify the particle antiparticle states of a given total angular momentum j with respect to P- and CP-parity:

(1)	$s=0,$	$l=j,$	$P=(-)^{j+1},$	$C=(-)^j,$	$CP=-$
(2)	$s=1,$	$l=j,$	$P=(-)^{j+1},$	$C=(-)^{j+1},$	$CP=+$
(3)	$s=1,$	$l=j+1,$	$P=(-)^j,$	$C=(-)^j,$	$CP=+$
(4)	$s=1,$	$l=j-1,$	$P=(-)^j,$	$C=(-)^j,$	$CP=+$

There are obviously only two CP-odd and P-even amplitudes:

$$1\rightarrow 2, 2\rightarrow 1.$$

By the way, the number of P-odd amplitudes, both CP-even and CP-odd, is larger, four in both cases (see, e.g., ref. [19]). Still larger, six, is the number of CP- and P-even amplitudes.

To construct explicitly TOPE amplitudes we recall that the T-odd, P-odd interaction of an electric dipole moment d with an electromagnetic field strength $F_{\mu\nu}$ is [19]

$$V_d = \frac{1}{2} d\bar{\psi}\gamma_5\sigma_{\mu\nu}\psi F_{\mu\nu} \tag{1}$$

Substituting a fermion vector current for the vector potential, we would get a four-fermion T-odd, P-odd interaction. Obviously, to obtain a T-odd, P-even one, we have to substitute for the vector potential an axial current. In this way we get two following four-fermion operators:

$$(G/\sqrt{2})(q_1/2m_p)\bar{\psi}_1 i\gamma_5\sigma_{\mu\nu}(p_1'-p_1)_\nu\psi_1\bar{\psi}_2\gamma_\mu\gamma_5\psi_2, \tag{2}$$

$$(G/\sqrt{2})(q_2/2m_p)\bar{\psi}_1\gamma_\mu\gamma_5\psi_1\bar{\psi}_2 i\gamma_5\sigma_{\mu\nu}(p_2'-p_2)_\nu\psi_2. \tag{3}$$

As was mentioned above, we measure the interaction discussed in the units of the Fermi weak interaction constant G; m_p is the proton mass, its choice as the necessary dimensional parameter being also a matter of convention; $q_{1,2}$ are dimensionless.

Let us note that these covariant operators necessarily contain derivatives. In other words, their dimension, seven, is higher than the lowest possible dimension, six, of other types of four-fermion operators.

The amplitudes (2), (3) could arise through the exchange by a neutral pseudovector boson, if its vertices contain the mixture of the "normal" axial operator $\gamma_\mu\gamma_5$ and the "anomalous" one $i\gamma_5\sigma_{\mu\nu}(p'-p)_\nu$ of opposite CP-parity [20].

On mass shell the above spinor structure can be transformed into

$$i\gamma_5(p'_1+p_1)_\mu \times \gamma_\mu\gamma_5, \tag{4}$$

$$\gamma_\mu\gamma_5 \times i\gamma_5(p'_2+p_2)_\mu. \tag{5}$$

We omit here evident scalar factors and field operators $\psi_{1,2}$; the first and second spin-momentum operators in both expressions refer to the first and second particles respectively. The representation (4), (5) is more convenient for most calculations.

The case of identical fermions ($\psi_1=\psi_2$) deserves special attention. Here the TOPE interaction vanishes if q's are independent of the momentum transfer. Technically, the direct and exchange amplitudes cancel after Fierz transformation. The following identities are useful for this proof and calculations below:

$$\sigma_{\mu\nu} \times \gamma_\mu(p'_2+p_2)_\nu = \gamma_5(p'_1+p_1)_\mu \times \gamma_\mu\gamma_5, \tag{6}$$

$$\sigma_{\mu\nu} \times \gamma_\mu\gamma_5(p'_2+p_2)_\nu = -m_2\gamma_5\sigma_{\mu\nu} \times \sigma_{\mu\nu} + \gamma_5\sigma_{\mu\nu}(p'_1-p_1)_\nu \times \gamma_\mu. \tag{7}$$

If one takes into account the possible dependence of q's on the invariant momentum transfers t and u in the direct and exchange channels, respectively, the TOPE interaction of identical fermions does not vanish anymore. But being proportional to, at least, $t-u$, this interaction is described by an operator of the dimension nine or higher.

T-ODD, P-EVEN ELECTROMAGNETIC INTERACTIONS

Let us start from a possible TOPE part of a one-photon vertex. The matrix element J_μ of the electromagnetic current operator between states with a given spin I and momenta k and k' can be conveniently expanded in the four independent Lorentz vectors of the problem:

$$p_\mu = (k'+k)_\mu,\ q_\mu = (k'-k)_\mu,\ s_\mu,\ r_\mu = i\epsilon_{\mu\nu\kappa\lambda}p_\nu q_\kappa s_\lambda. \tag{8}$$

Here s_μ is the four-dimensional spin operator defined, e.g., for the state with momentum k. It is obtained through the Lorentz transformation of the spin vector $(0,\mathbf{I})$ in the rest frame of a particle and can be written as

$$s_\mu = (s_0,\mathbf{s}),\ s_0 = (\mathbf{I}\cdot\mathbf{k})/m,\ \mathbf{s} = \mathbf{I} + (\mathbf{I}\cdot\mathbf{k})\mathbf{k}/m(E+m).$$

Being defined in this way, the vector s_μ obviously satisfies the condition $k_\mu s_\mu = 0$ (in which case, however, $k'_\mu s_\mu \neq 0$). Allowing for the current conservation law which is in the momentum representation $q_\mu J_\mu = 0$, the matrix element of the current can be written as follows (see, e.g., ref. [19]):

$$J_\mu = \langle \mathbf{I},\mathbf{k}'|p_\mu F_1 + r_\mu F_2 + [q^2 s_\mu - q_\mu(qs)]F_3|\mathbf{I},\mathbf{k}\rangle\ . \tag{9}$$

The invariant functions F_i depend on the scalars that can be constructed from the vectors (6). Since pq, rp, rq, rs and sk vanish, there are only two such independent invariants. It is convenient to choose as arguments of F_i the invariant momentum transfer $t = q^2$ and the hermitian operator $\tau = iq_\mu s_\mu$ which is P-odd, T-odd and C-even. The substitution of the operator function $F_i(t,\tau)$ into expression (9) should be accompanied by symmetrization in the non-commuting operators s_μ and τ. In the expansion of $F_i(t,\tau)$ in powers of τ

$$F_i(t,\tau) = \sum_{n=0}^{N_i} f_{in}(t)\tau^n \tag{10}$$

the highest power N_i is evidently fixed by the spin I:

$$N_1 = 2I, N_2 = N_3 = 2I - 1.$$

Clearly, the first two structures in the matrix element (9) possess the correct C-parity, odd terms of the expansion of $F_{1,2}$ violating simultaneously P and T. But the axial vector $q^2 s_\mu - q_\mu(qs)$ behaves properly under T-reversal and is therefore of wrong C-parity. The electromagnetic interaction of J_μ with the test current j_μ is written as

$$-j_\mu \frac{1}{q^2} J_\mu. \tag{11}$$

Since the test current j_μ is conserved, $q_\mu s_\mu = 0$, the term $-q_\mu(qs)$ in J_μ is not operative, while the factor q^2 in $q^2 s_\mu$ cancels the propagator $1/q^2$ in (8). The remaining C-odd interaction $-j_\mu < I, k'|s_\mu F_3|I, k >$ is a local one. To summarize, in the case of spin 1/2 there is no TOPE one-photon electromagnetic interaction. And for an arbitrary spin any C-odd electromagnetic vertex, both P-odd and P-even, is necessarily of a contact nature. This conclusion has been made many years ago in ref. [21].

Now we are going over to TOPE electromagnetic amplitudes of photon-fermion scattering, to two-photon interaction, or Compton one. A standard analysis (see, e.g., ref. [22]) demonstrates that there are two such amplitudes, nonvanishing for real photons. Below we will present a derivation which not only allows one to find the number of those amplitudes, but to derive them in an explicitly gauge-invariant form. This form seems to be more convenient, at least, for our purposes, than the standard one presented in ref. [22].

Again, it is convenient to start from the annihilation channel $f\bar{f} \to 2\gamma$. The C-parity of the 2γ state is certainly positive. And since the interaction we are interested in, violates the invariance under charge conjugation, the C-parity of the fermion-antifermion state should be negative.

At the vanishing total spin of the fermionic pair, $S = 0$, this means that its orbital angular momentum l should be odd and the parity positive, $P = +$. But at $l = 1$ the total angular momentum of a singlet state is $J = 1$ which is impossible for two photons (see, e.g., ref. [22]). Thus, one possibility corresponds to the annihilation in the positive-parity singlet states with odd total angular momenta starting from $J = 3$.

At $S = 1$ the negative C-parity of the fermion-antifermion pair implies even l and, correspondingly, $P = -$. And again $l = 0$ leads at $S = 1$ to $J = 1$ which is forbidden. Since the two-photon states with odd J have necessarily positive parity, $P = +$ [22], the second possibility corresponds to the annihilation from the negative-parity triplet state with even total angular momenta starting from $J = 2$.

In the last case the amplitude is more simple and looks as follows:

$$\gamma_5[\gamma_\mu(p'-p)_\nu + \gamma_\nu(p'-p)_\mu](k'-k)_\mu(k'-k)_\nu F_{\alpha\beta}\tilde{F}_{\alpha\beta}. \tag{12}$$

Here p and p' are the momenta of the annihilating electron and positron respectively, k and k' are those of the photons, $\tilde{F}_{\alpha\beta} = \frac{1}{2}\epsilon_{\alpha\beta\gamma\delta}F_{\gamma\delta}$. To go over to the scattering channel we have to change p' to $-p'$, and k' to $-k$. The interaction discussed is of a high dimension, ten, not only due to two photons participating, but first of all due to a large number of momenta, or derivatives, in it. The reason is that the angular momenta in the annihilation channel are high, starting from $J = 2$. Finally, the interaction discussed can be presented as

$$\frac{A}{8m_p^6}(\bar{\psi}\gamma_\mu\gamma_5 i\overleftrightarrow{\partial}_\nu\psi)(F_{\alpha\beta}i\overleftrightarrow{\partial}_\mu i\overleftrightarrow{\partial}_\nu\tilde{F}_{\alpha\beta}). \tag{13}$$

We single out here a dimensionless constant A, taking again as the necessary dimensional parameter the proton mass m_p. The symbol $\overleftrightarrow{\partial}$ is $\overleftrightarrow{\partial} = \vec{\partial} - \overleftarrow{\partial}$ where the derivatives $\vec{\partial}$ and $\overleftarrow{\partial}$ are acting to the right and to the left, respectively. The annihilation from higher even angular momenta can be described by making A momentum-transfer dependent.

As to the amplitude corresponding to the annihilation from the singlet state and starting from $J = 3$, its fermionic part should be evidently

$$i\gamma_5 P_\kappa P_\lambda P_\mu;\ P_\kappa = (p'-p)_\kappa.$$

Then one should guarantee the Bose-statistics, the symmetry of the interaction under the permutation of the photons in the annihilation channel. In this way we come to the following form of the operator discussed:

$$\frac{B}{16m_p^9}(\bar{\psi}i\gamma_5 i\overleftrightarrow{\partial}_\kappa i\overleftrightarrow{\partial}_\lambda i\overleftrightarrow{\partial}_\mu\psi)(i\partial_\rho F_{\kappa\sigma}i\overleftrightarrow{\partial}_\mu i\partial_\sigma \tilde{F}_{\lambda\rho}). \tag{14}$$

The dimension of this operator, thirteen, is even higher than that of the previous one. That is why we have to introduce the factor m_p^{-9} at the dimensionless constant B.

Now one can easily obtain possible three-photon TOPE fermion amplitudes. It is achieved by introducing into (13), (14) the extra "factor" $\sigma_{\gamma\delta}F_{\gamma\delta}$. Then one has to single out independent spinorial structures and to introduce, if necessary, the factor i to make the operator hermitian. In this way we get

$$[\bar{\psi}\gamma_5(\gamma_\mu i\overleftrightarrow{\partial}_\nu + \gamma_\nu i\overleftrightarrow{\partial}_\mu)\psi]F_{\mu\lambda}(F_{\rho\sigma}i\overleftrightarrow{\partial}_\lambda i\overleftrightarrow{\partial}_\nu\tilde{F}_{\rho\sigma}), \tag{15}$$

$$[\bar{\psi}(\gamma_\mu i\overleftrightarrow{\partial}_\nu + \gamma_\nu i\overleftrightarrow{\partial}_\mu)\psi]\tilde{F}_{\mu\lambda}(F_{\rho\sigma}i\overleftrightarrow{\partial}_\lambda i\overleftrightarrow{\partial}_\nu\tilde{F}_{\rho\sigma}), \tag{16}$$

$$(\bar{\psi}i\sigma_{\alpha\beta}i\overleftrightarrow{\partial}_\kappa i\overleftrightarrow{\partial}_\lambda i\overleftrightarrow{\partial}_\mu\psi)\tilde{F}_{\alpha\beta}[(i\partial_\rho F_{\kappa\sigma})i\overleftrightarrow{\partial}_\mu(i\partial_\sigma\tilde{F}_{\lambda\rho})]. \tag{17}$$

A technical remark: one can get rid of tildes in (16), (17), but then the expressions become less compact.

Interactions (15) and (16) describe TOPE amplitudes of the three-photon annihilation of a fermion-antifermion pair with even total angular momenta starting from $J = 2$ and of negative and positive parity, respectively. Expression (17) should be expanded into irreducible structures. It corresponds to the annihilation from the positive-parity states of any J starting from 2^+.

But what about the three-photon annihilation from the fermion-antifermion ground states ${}^1S_0$ and ${}^3S_1$? The triplet state is of the same negative C-parity, as the three-photon one, so for it there is no TOPE amplitude at all. The P-even interaction which could be responsible for the three-photon decay of parapositronium or π^0-meson, was pointed out long ago in ref. [23]. In our notations it is

$$(\bar{\psi} i\gamma_5 \psi)\partial_\alpha \tilde{F}_{\rho\sigma}\partial_\beta\partial_\gamma F_{\rho\sigma}\partial_\gamma F_{\alpha\beta}. \tag{18}$$

One more operator of the type discussed can be obtained from the anapole interaction by multiplying it by $F\tilde{F}$. But in this case one of the three photons cannot be real, and we will not consider here the interactions of this kind.

NEW LIMITS ON T-ODD, P-EVEN FERMION-FERMION INTERACTIONS. 1

After this purely kinematical discussion let us go over to the calculation of the electroweak radiative corrections which transform T-odd, P-even operators (2), (3) into T- and P-odd ones. A hint at the kind of limits that can be obtained in this way is given by the following argument, close in spirit to the corresponding estimates from refs. [24,25]. Let us consider the contribution to the neutron electric dipole moment from the combined action of the usual P-odd, T-even weak interaction and the discussed T-odd and P-even interaction, the strength of the latter being q times smaller than that of the previous one. The contribution constitutes obviously

$$d_n/e \sim m_p^{-1}(Gm_\pi^2)^2 q\, cm \tag{19}$$

From the comparison with the last experimental results [26, 27]

$$d_n/e < 10^{-25}\, cm \tag{20}$$

we get the limit $q < 10^2$, which is about two orders of magnitude better than the direct limits mentioned in the Introduction. This estimate is obviously of a very crude nature. In particular, the dipole moment arises here at least in one-loop approximation which leads to a small geometrical factor. It suppresses the above estimate at least by an order of magnitude. So, in this way we cannot get the limit better than

$$q < 10^3. \tag{21}$$

Still, this example is quite instructive.

The first improvement of the above limit is due to the observation that the electroweak corrections to the operators (2), (3) are controlled mainly not by the large-distance effects, but by the short-distance ones. Therefore, they are of the order α/π (up to some chiral suppression factor which is quite essential), but not of the order Gm_π^2.

We will concentrate here and below on the corrections due to the Z-boson exchange mainly. They can be calculated self-consistently in the sense that the result is independent of the choice of the gauge for the Z-boson propagator. The result is independent also of the form of the initial T-odd operators (2), (3) or (4), (5).

A consistent, gauge-independent calculation of the W-boson exchange contribution to the induced T- and P-odd amplitudes is much more model-dependent and will not be discussed here in detail. It can be expected however to be even larger than that of the Z-exchange, due to small numerical values of the neutral weak charges. These small values are responsible in particular for the well-known relative suppression of the neutral-current cross-sections as compared to the charged-current ones. So, the Z-exchange contribution serves as an estimate from below for the effects discussed. As to the Higgs boson exchange, in the standard model it conserves parity and is therefore of no interest to us.

The details of the calculations can be found in ref. [16] and here I will present just the result of the transformation of TOPE operator (2) into T- and P-odd one, which is

$$\frac{G}{\sqrt{2}}\frac{\alpha}{3\pi}\log\frac{\Lambda^2}{M^2}(q_1/m_p)\{2m_2v_1[3a_2i\gamma_5\times 1-(a_1+a_2)1\times\gamma_5] \\ -a_1v_2[m_2i\gamma_5\sigma_{\mu\nu}\times\sigma_{\mu\nu}-i\gamma_5\sigma_{\mu\nu}(p_1'-p_1)_\nu\times\gamma_\mu]\} \tag{22}$$

Here M is the Z-boson mass. The dependence of the result on the cut-off parameter Λ is due to nonrenormalizability of the TOPE interaction. But trying to be as conservative as possible in our numerical estimates, we will assume the *log* to be of the order of unity. $m_{1,2}$ are the masses of the first and second fermions respectively. $v_{1,2}$ and $a_{1,2}$ are their weak neutral vector and axial charges. In particular, for the electron, u-, and d-quarks they are:

$$\begin{aligned} v_e &= -\frac{1}{2}(1-4\sin^2\theta)\approx -0.04, \quad a_e=-\tfrac{1}{2}, \\ v_u &= \frac{1}{2}(1-\frac{8}{3}\sin^2\theta)\approx\frac{1}{6}, \quad a_u=\tfrac{1}{2}, \\ v_d &= -\frac{1}{2}(1-\frac{4}{3}\sin^2\theta)\approx-\frac{1}{3}, \quad a_d=-\tfrac{1}{2}, \\ &\sin^2\theta\approx 0.23. \end{aligned} \tag{23}$$

Let us perform the concrete estimates for the electron-nucleon interaction. In this case the induced electron-quark operator is

$$\frac{G}{\sqrt{2}}\frac{\alpha}{3\pi}\log\frac{\Lambda^2}{M^2}v\{q_e[\frac{m}{m_p}i\gamma_5\sigma_{\mu\nu}\times\sigma_{\mu\nu}-\frac{1}{2m_p}i\gamma_5\sigma_{\mu\nu}(p_1'-p_1)_\nu\times\gamma_\mu] \\ +q\frac{m_e}{m_p}[(1-2a)i\gamma_5\times 1-3\cdot 1\times i\gamma_5]\}. \tag{24}$$

Here m and m_e are the quark and electron masses respectively, v and a are the quark vector and axial charges, q_e and q are the dimensionless constants in the T-odd, P-even operators with the explicit momenta belonging to electrons and quarks, respectively. In all the products the first operator refers to electrons, the second one to quarks. We neglect here the contribution proportional to the electron vector charge v_e which is numerically small (see (23)). Operator (24) should be summed over u- and d-quarks, and its expectation value should be taken first over a nucleon and then over a nucleus.

In the static approximation for nucleons, the only term in (24) that depends on both electron and nucleon spin is $\gamma_5\sigma_{\mu\nu} \times \sigma_{\mu\nu}$. The dimensional estimate for the nucleon expectation value of the quark operator $\sigma_{\mu\nu}$ is $\bar{N}\sigma_{\mu\nu}N$. Then the dimensionless effective constant of the T- and P-odd tensor electron-neutron interaction is

$$k_2 = \frac{\alpha}{3\pi}\log\frac{\Lambda^2}{M^2}(m_d/2m_p)v_d q_e \sim 10^{-6}q_e. \qquad (25)$$

The upper limit on the constant k_2 that follows from the experiment with ^{199}Hg [28] leads to the following result:

$$q_e < 1. \qquad (26)$$

For the d-quark mass we assume here the value $m_d = 7\,MeV$. Close limit on q_e follows from the experiment with TlF [29].

Almost the same upper limits can be extracted in this way from atomic experiments, for the constants $q_{u,d}$ referring to the electron-quark interaction with the derivatives in the quark vertex.

The analogous estimates for the quark-quark constants (see ref. [16]) are more tedious, due mainly to the problem of evaluating the contributions of the arising T- and P-odd four-quark operators to the neutron EDM. In this way one gets for different quark-quark constants the upper limits on the level

$$q < 10. \qquad (27)$$

We do not go here into details since in the next section much better upper limits for all these constants will be obtained in the two-loop approximation.

NEW LIMITS ON T-ODD, P-EVEN FERMION-FERMION INTERACTIONS. 2

The idea of the next improvement of the upper limits on the constants discussed, can be conveniently explained for the case of hadrons. The previous improvement from (21) to (27) has been reached by going over from long-distance effects of the usual weak interaction to the short-distance ones. It allowed us to get rid of one small factor Gm_π^2 in formula (19), trading it for α/π with some extra chiral suppression. But cannot we get rid of the second factor Gm_π^2 in that formula? The answer is: yes, we can. Up to now we computed one-loop radiative correction which transformed the interaction discussed into T- and P-odd effective operator. Then we estimated in fact the long-distance contribution of this operator to the neutron electric dipole moment (EDM). Now we are going to make the next step: to calculate a completely short-distance two-loop contribution of the TOPE interaction times the weak interaction, directly to the quark dipole EDM. And the latter is of the same order of magnitude as the neutron dipole moment. The gain is even larger than that at the first step since now there is no more chiral suppression factor.

To regularize, at least, partly, Feynman integrals, it is convenient to introduce explicitly the axial boson of mass μ mediating the TOPE interaction. Then the TOPE fermion-fermion amplitude can be presented as

$$\frac{4\pi\beta}{\mu^2-k^2}\frac{1}{2m_p}\bar{\psi}_1 i\gamma_5\sigma_{\mu\nu}(p_1'-p_1)_\nu\psi_1\bar{\psi}_2\gamma_\mu\gamma_5\psi_2. \tag{28}$$

Here β is the dimensionless coupling constant analogous to $\alpha = 1/137$ in QED.

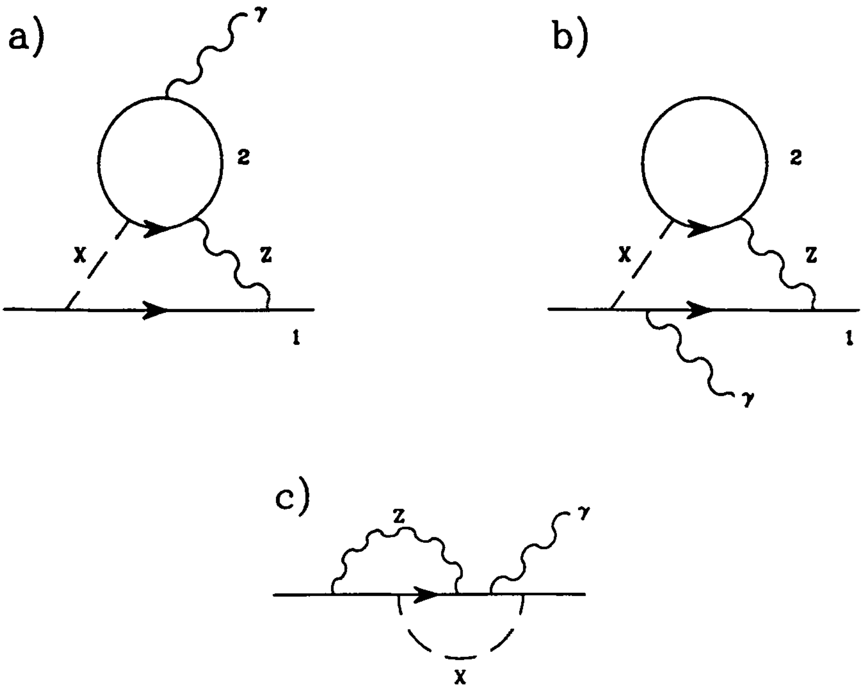

Figure 1: Two-loop diagrams

Consider now the contribution of the two-loop diagram 1a to the fermion EDM. Here the dashed line represents the propagator of the axial boson X. The Z-boson exchange (the wavy line) introduces the parity nonconservation necessary to induce the EDM which is both T- and P-odd.

It can be easily seen that the contributions proportional to the fermion masses are as small as the Fermi constant G. Therefore, to our accuracy the fermions will be taken as massless. Since the dipole moment interaction $\frac{1}{2}\gamma_5\sigma_{\alpha\beta}F_{\alpha\beta}$ changes the fermion chirality, the lower vertex of the X-line should be $i\gamma_5\sigma_{\mu\lambda}k_\lambda$ and the upper one, correspondingly, $\gamma_\mu\gamma_5$. Then according to the Furry theorem for the fermion loop, the upper vertex of the Z line is γ_ν, and the lower one, respectively, $\gamma_\nu\gamma_5$. The fermion loop arising in this way has in fact been calculated in ref. [30]. For vanishing fermion mass the result is

$$\frac{i}{8\pi^2}F_{\alpha\beta}[\epsilon_{\alpha\beta\mu\nu} + \frac{1}{k^2}\epsilon_{\alpha\beta\kappa\lambda}k_\kappa(\delta_{\mu\lambda}k_\nu + \delta_{\nu\lambda}k_\mu)]; \tag{29}$$

here $F_{\alpha\beta}$ is the strength of the external field.

When considering the lower, Compton block of diagram 1*a*, one should also include the contact term

$$\frac{e}{2\sin\theta\cos\theta}2ai\sigma_{\mu\nu} \tag{30}$$

originating from the vertex $i\gamma_5\sigma_{\mu\lambda}(p'-p)_\lambda$ via the substitution

$$p_\mu\psi \to [p_\mu - \frac{e}{2\sin\theta\cos\theta}(v + a\gamma_5)Z_\mu]\psi \tag{31}$$

which makes this vertex covariant with respect to the Z-field. Here e is the electric charge, θ is the electroweak mixing angle, v and a are the vector and axial charges, respectively. In particular, the inclusion of the contact vertex (30) into the XZ Compton scattering amplitude makes the result of the calculation independent of the term $k_\nu k_\sigma/M^2$ in the propagator of the Z-boson. Simple calculations give the following result for this contribution to the EDM d of the fermion 1:

$$\frac{d}{e} = -\frac{\alpha\beta Q_2 a_1 v_2}{3\pi^2 m}\log\frac{\Lambda^2}{M_>^2}. \tag{32}$$

This formula refers to the general case in which the fermion 2 propagating in the loop differs from the fermion 1 propagating in the lower line. In particular, a_1 is the axial weak charge of the first fermion, v_2 is the vector weak charge of the second fermion, and Q_2 is its electric charge in the units of e. The logarithmic dependence on the cutoff parameter Λ is the result of the nonrenormalizable coupling of the X-boson to the vertex with derivatives $i\gamma_5\sigma_{\mu\rho}k_\rho$. Although the result (32) is presented with logarithmic accuracy, in all of our numerical estimates we will conservatively assume $\log\Lambda^2/M_>^2$ to be of the order of unity ($M_>$ is the largest of the masses μ and M).

The contribution of diagram 1b

$$\frac{d_2}{e} = \frac{\alpha\beta Q_1 a_2 v_1}{36\pi^2 m}\log^2\frac{\Lambda^2}{M_>^2} \tag{33}$$

is much smaller numerically and can be neglected.

In the case of identical fermions there are also the contributions to the EDM of diagrams of type 1c, but we will neglect them in our estimates with the expectation that the results will not be grossly affected. In particular, even in the case $\mu \gg M$ we can, to logarithmic accuracy, restrict the integration over k to $k \gg \mu$, where the TOPE interaction of identical fermions is in no way a local one and therefore should not vanish.

A consistent, gauge-independent calculation of the W-boson exchange contribution to the induced EDM is again much more model-dependent and will not be discussed in this paper. The same arguments as in the previous section, allow us to expect that the Z-boson contribution serves as a conservative estimate for the induced EDM.

We now consider the two-loop contribution to the electron EDM due to the electron-electron TOPE interaction. Substituting into formula (32) the numerical values (23) for a_e, v_e, as well as $Q_e = -1$, we get

$$\frac{d_e}{e} \sim \beta_e \cdot 10^{-19}\, cm. \tag{34}$$

The upper limit on the electron EDM that follows from the atomic experiments [31,32]

$$\frac{d_e}{e} < 10^{-26}\, cm, \tag{35}$$

leads to the following result for the constant β_e of the electron-electron TOPE interaction:

$$\beta_e < 10^{-7}. \tag{36}$$

In the same way we can get new, very strict upper limits on the electron-nucleon and nucleon-nucleon TOPE interactions. The axial charge of a fermion is always (up to a sign) 1/2 and for any quark, independently of its sort, the product Qv is numerically close to 1/9. Then, using the experimental upper limit (20) on the neutron EDM and assuming for dimensional reasons that the neutron dipole moment induced by the quark EDM has about the same magnitude as the latter, we get for β_{qe}, the TOPE quark-electron interaction constant with derivatives in the quark vertex, the limit

$$\beta_{qe} < 10^{-6}. \tag{37}$$

For another distinct electron-quark constant β_{eq} (with the derivative in the electron vertex) the constraint (35) on the electron EDM gives the upper limit

$$\beta_{eq} < 3 \times 10^{-8}. \tag{38}$$

For all quark-quark constants β_{qq} the limit (20) on the neutron EDM gives

$$\beta_{qq} < 3 \times 10^{-7}. \tag{39}$$

The latter refers as well to the "coloured" TOPE interaction, with the SU(3) generators t^a in each vertex. In this case the external field on diagram 1 should not be electromagnetic, but rather a gluon one. Again, for dimensional reasons, the neutron EDM induced in this way should be of the same order of magnitude as the colour dipole moment described by this diagram.

The constants β introduced here, are related as follows to the constants q used above:

$$\frac{4\pi\beta}{\mu^2} = \frac{G}{\sqrt{2}}q, \tag{40}$$

or

$$q = 4\pi\beta\sqrt{2}(m_p/\mu)^2 \times 10^5 = 1.8 \times 10^6 (m_p/\mu)^2 \beta. \tag{41}$$

Thus, the upper limits corresponding to (37), (38) and (39) are

$$q_{qe} < 2(m_p/\mu)^2,\ q_{eq} < 0.05(m_p/\mu)^2,\ q_{qq} < 0.5(m_p/\mu)^2, \tag{42}$$

respectively.

The upper limits on the constants q, close to $q < 10$, were derived in the previous section under the assumption $\mu \geq M \sim 100m_p$. Under the same assumption about μ, the limits we obtain here are much better:

$$q_{qe} < 10^{-4},\ q_{eq} < 10^{-5},\ q_{qq} < 10^{-4}. \tag{43}$$

Let us come back to the explanation of this gain. In the transition from the effective four-fermion T- and P-odd operators obtained in the previous section to the neutron EDM we used the usual hadronic scale of 1 GeV. But here the transition takes place on a much higher scale of 100 GeV.

ELECTRON-NUCLEON AND NUCLEON-NUCLEON INTERACTIONS

Now, having obtained the above limits on the TOPE electron-quark and quark-quark interactions, what can we say about the corresponding electron-nucleon and nucleon-nucleon interactions?

The answer for the electron-nucleon interaction is a quite straightforward one. Simple dimensional arguments lead to the following estimates for the nucleon expectation values of the relevant quark operators (the last of them has been already mentioned):

$$< N|\bar{q}\gamma_\mu\gamma_5 q|N >\sim \bar{N}\gamma_\mu\gamma_5 N,\ < N|\bar{q}\gamma_5\sigma_{\mu\nu}q|N >\sim \bar{N}\gamma_5\sigma_{\mu\nu}N. \tag{44}$$

Therefore, the limits (42) for $q_{qe,eq}$ are readily translated into those for the constants of TOPE electron-nucleon interactions:

$$q_{Ne} < 10^{-4}, q_{eN} < 10^{-5}. \tag{45}$$

Let us address now the nucleon-nucleon interactions. Note first of all that in contrast to T- and P-odd nuclear forces, TOPE ones cannot be mediated by π^0-meson exchange. Indeed, looking at the classification of the particle-antiparticle states in the annihilation channel, we see that at $j = 0$ the state 2 just does not exist.

The absence of this exchange can be attributed also to vanishing of a TOPE $\pi^0 NN$ vertex. As to the TOPE $\pi^\pm NN$ coupling, being hermitian it looks like

$$\bar{p}\gamma_5 n\pi^+ - \bar{n}\gamma_5 p\pi^-. \tag{46}$$

This coupling does not lead to TOPE NN scattering amplitude in the one-boson exchange approximation since after the interchange of this vertex and of the strong one, the corresponding diagrams cancel out. Vanishing of the one-pion exchange in the TOPE NN scattering was pointed out previously in ref. [33]. TOPE one-boson exchange starts therefore with vector and pseudovector bosons. Being mediated by heavier particles, the effective NN interaction is suppressed as compared to simple estimates.

On the other hand, it follows already from general formulae (2) and (3) that a TOPE nucleon-nucleon scattering amplitude contains an extra power of p/m_p as compared to the usual P-odd weak interaction. This means a suppression of roughly by an order of magnitude as compared with the mentioned naive estimate $Gm_\pi^2 q$.

Thus, even taking into account all the uncertainties of our estimates, one can expect that the relative strength of the TOPE nuclear forces does not exceed $10^{-4}Gm_\pi^2$.

One should have in mind however that the observable effects of the TOPE nuclear forces may exceed this estimate, due to long-distance enhancement factors, such as small energy intervals between resonances mixed by the interaction discussed.

In the conclusion of this section let us note that in the same way one can obtain upper limits on the constants of the photon-fermion interactions.

β-DECAY CONSTANTS

Some information can be obtained in an analogous way concerning even β-decay constants. To relate them to the eN TOPE interaction one should evidently switch on W-exchange. As has been mentioned already, this procedure is more ambiguous than the switching on of Z-exchange used in our previous consideration. One can hope, however, that the estimates made below are valid at least by an order of magnitude.

Let us start from the consideration of T-odd quark-lepton β-decay interaction without derivatives. Then, to obtain the radiative correction of the order of magnitude α/π, but not G, the quark mass should be neglected. Let us note that all T- and P-odd fermion-fermion interaction operators without derivatives change the chirality of both fermions. Since the W-exchange vertices contain left projectors, we can investigate in this way, neglecting the quark masses, the chirality-changing quark-lepton operators only:

$$(\bar{u}d)[\bar{e}(c_S + c'_S\gamma_5)\nu] + (\bar{u}\gamma_5 d)[\bar{e}(c_P\gamma_5 + c'_P)\nu] - \frac{1}{2}(\bar{u}\sigma_{\mu\nu}d)[\bar{e}(c_T + c'_T\gamma_5)\sigma_{\mu\nu}\nu] + h.c. \quad (47)$$

So, in this approximation nothing can be said about axial and vector constants. On the other hand, the P-even part of the electroweak correction allows us to obtain information on the P- and T-odd β-decay constants as well.

The effective T- and P-odd interaction of electron with u-quark arising from (47) through W-exchange is

$$\frac{G}{\sqrt{2}}\frac{\alpha}{4\pi}\log\frac{\Lambda^2}{M_W^2}Im[6(c_T+c'_T)-(c_S+c'_S)-(c_P+c'_P)][(i\gamma_5\times 1+1\times\gamma_5)+\frac{i}{2}\gamma_5\sigma_{\mu\nu}\times\sigma_{\mu\nu}]. \quad (48)$$

The corresponding effective operator for the interaction between electron and d-quark is

$$\frac{G}{\sqrt{2}}\frac{\alpha}{4\pi}\log\frac{\Lambda^2}{M_W^2}Im[(c_S + c'_S) - (c_P + c'_P)][i\gamma_5 \times 1 - 1 \times i\gamma_5]. \quad (49)$$

Again in each operator product in these formulae the first factor refers to the electron, the second one refers to the quark. In all our calculations we assume for the W-boson propagator the simple Feynman form $\delta_{\mu\nu}/(q^2 - M_W^2)$ with the hope that the term $-q_\mu q_\nu/M_W^2$ in the numerator will somehow cancel out at more accurate calculations.

Now we get for the effective constant k_2 (see the discussion of formulae (25), (26)) expression

$$k_2 = \frac{\alpha}{8\pi} \log \frac{\Lambda^2}{M_W^2} Im[6(c_T + c'_T) - (c_S + c'_S) - (c_P + c'_P)], \tag{50}$$

and the following limit on the β-decay constants:

$$Im[6(c_T + c'_T) - (c_S + c'_S) - (c_P + c'_P)] < 4 \times 10^{-3} \tag{51}$$

Let us turn now to the eN interaction independent of nuclear spin, which is generated by the structure $i\gamma_5 \times 1$ in formulae (48), (49). The proton and neutron expectation values for the quark scalar operators equal [34-36]

$$< p|\bar{u}u|p >=< n|\bar{d}d|n >= 6, \ < p|\bar{d}d|p >=< n|\bar{u}u|n >= 5. \tag{52}$$

So, the expectation values of those operators over a heavy nucleus are

$$< \bar{u}u > \approx < \bar{d}d > \approx 5.5A \tag{53}$$

where A is the atomic number. In this way we get for the induced T- and P-odd electron-nucleus interaction constant the following expression:

$$k_1 = \frac{\alpha}{2\pi} \log \frac{\Lambda^2}{M_W^2} 5.5 Im[3(c_T + c'_T) - (c_P + c'_P)]. \tag{54}$$

The experiment [32] leads to the limit on the constant k_1 on the level 2.3×10^{-6} (see ref.[19]). In this way we come to the constraint for another combination of T-odd β-decay constants:

$$Im[3(c_T + c'_T) - (c_P + c'_P)] < 4 \times 10^{-4} \tag{55}$$

So, the final results for the T-odd quark-lepton β-decay constants can be formulated (up to the possibility of some cancellations) as

$$Im(c_S + c'_S) < 4 \times 10^{-3}, \ Im(c_T + c'_T) < 10^{-4}, \ Im(c_P + c'_P) < 4 \times 10^{-4}. \tag{56}$$

Finally we have to pass over from the quark β-decay constants to the usual ones that refer to nucleons. The scalar and pseudoscalar matrix elements are

$$< p|\bar{u}d|n >= \frac{m_\Xi - m_\Sigma}{m_s} \bar{p}n, \ < p|\bar{u}i\gamma_5 d|n >= \frac{f_\pi g_r \sqrt{2}}{m_u + m_d} \bar{p}i\gamma_5 n \tag{57}$$

where $m_{\Xi,\Sigma}$ are the hyperon masses, $m_s = 140$ MeV and $m_u = 4$ MeV are the masses of the s- and u-quarks, $f_\pi = 130$ Mev, $g_r = 13.5$ is the renormalized strong coupling constant. For the tensor matrix element we use an order of magnitude estimate

$$< p|\bar{u}\sigma_{\mu\nu}d|n > \sim \bar{p}\sigma_{\mu\nu}n. \tag{58}$$

In this way we come to the following upper limits on the nucleon β-decay constants:

$$Im(C_S + C_S') < 4 \times 10^{-3},\ Im(C_T + C_T') < 10^{-4},\ Im(C_P + C_P') < 0.1. \tag{59}$$

The limits on diverse combinations of Im$C_{S,T,P}$ obtained from direct experiments [10-12, 14,15] lie roughly between 10^{-2} and unity, depending on the combination involved. All of those combinations differ, however, from ours. But what is more essential, our limits depend on some assumptions and are of a less quantitative character. Nevertheless, I believe that the limits (59), being much more stringent than the direct ones, are quite interesting.

The transition to the two-loop approximation does not tell us anything interesting about the T-odd β-decay interaction without derivatives. But it does about the so-called weak magnetism and weak dipole moment, i.e., about the β-decay derivative coupling which is on the quark-lepton level

$$\frac{G}{\sqrt{2}}\frac{1}{2m_p}[\bar{e}\gamma_\mu(1+\gamma_5)\nu][\bar{u}(\tilde{g}_m + \tilde{g}_e\gamma_5)\sigma_{\mu\nu}k_\nu d] + h.c. \tag{60}$$

The constants $Im\tilde{g}_{m,e}$ describing T-odd part of this interaction are close analogues of the dimensionless parameters q introduced above. No wonder therefore that in the two-loop approximation we get here the same upper limits

$$Im\tilde{g}_{m,e} < 10^{-4}. \tag{61}$$

Dimensional considerations show that the same limits as (61) for the quark-lepton β-decay constants, hold also for the parameters of the nucleon weak current induced in this way:

$$Im g_{m,e} < 10^{-4}. \tag{62}$$

The recent experimental proposal [37] aims at the accuracy $\sim 10^{-2}$ in the measurement of these constants.

T-odd correlations in the nuclear β-decays as well may be considerably enhanced due to "long-distance" effects in a nucleus. In this connection I wish to mention the possibility of such an enhancement due to the presence of an anomalously close nuclear level of opposite parity admixed by the T- and P-odd nucleon-nucleon interaction to the initial or final level.

The last remark in this section refers to the strangeness-changing β-decay interactions which differ from (47) and (60) by the substitution of the s-quark for the d-quark. Its T-odd constants are also limited from above by the results of atomic and neutron experiments. The limits are roughly five times worse than (59), (62), due mainly to the presence of the Cabibbo angle in the strangeness-changing W-exchange. As to the chirality-changing strange quark operators $\bar{s}s, \bar{s}i\gamma_5 s$, and $\bar{s}\sigma_{\mu\nu}s$, their nucleon expectation values are of about the same magnitude as those for d-quarks [35,36]. Quite meaningful limits follow in this way even on the T-odd constants for β-decay of heavy quarks: $c \to de\nu, b \to ue\nu, t \to de\nu$. But here again "long-distance" effects, well-known in these cases, lead to a tremendous enhancement of CP-odd phenomena in the semileptonic decays of neutral mesons, observed long ago in the decays $K^0 \to \pi l\nu$.

NONRELATIVISTIC REDUCTION OF TOPE AMPLITUDES

In this, final section of the talk we will consider the nonrelativistic limit of TOPE fermion-fermion and photon-fermion interactions, this limit being convenient for the discussion of future atomic experiments aimed at the investigation of TOPE interactions (see ref. [8]).

Let us start from positronium. Here in the local limit ($\mu \to \infty, 4\pi\beta/\mu^2$ =const) the "scattering" term is cancelled by its "annihilation" counterpart. So, we shall address the opposite limiting case of μ small as compared to the typical atomic momenta $m_e\alpha$. Of course, this (pseudo)vector boson interacting with a nonconserved axial current $\bar{\psi}\gamma_\mu\gamma_5\psi$ cannot be massless. It is not trivial to accomodate a light axial boson in a self-consistent theory, but we will not consider this problem here, confining ourselves to the discussion of present and future experimental limits on β. In any case, the term $k_\mu k_\nu/\mu^2$, singular in μ, in the numerator of the boson propagator does not contribute to amplitude (28) due to the conservation of the current $\bar{\psi}i\gamma_5\sigma_{\mu\rho}k_\rho\psi$.

We neglect μ altogether and confine ourselves to interaction (28), omitting the annihilation contribution, suppressed here, as compared to the scattering contribution, by a factor $\sim (k/\mu_e)^2 \sim \alpha^2$. The TOPE nonrelativistic interaction operator obtained in this way is

$$V = -\frac{\beta_e}{m_e m_p r^3}[\sigma_1 \times \sigma_2] \cdot \mathbf{l}. \tag{63}$$

Here σ_1 and σ_2 are the spin operators of the electron and positron, respectively, and $\mathbf{l}$ is the orbital angular momentum operator. This T-odd perturbation admixes to each other only the singlet and triplet states of the equal orbital angular momenta. These are the only positronium states of the same total angular momentum and parity with opposite C-parity. The mixing matrix element equals

$$< n^1l_l|V|n^3l_l >= -i\frac{\alpha\beta_e m_e}{2m_p}Ry\frac{1}{n^3(l+1/2)\sqrt{l(l+1)}} \tag{64}$$

where the Rydberg constant $Ry = 3.3 \times 10^{15}$Hz and n is the principal quantum number.

The eventual precision expected for the measurement of this matrix element between the states 2^1P_1 and 2^3P_1 in positronium is $\sim 10^5$Hz [9]. Therefore, the expected upper limit on the constant β_e is

$$\beta_e < 3 \times 10^{-4}. \tag{65}$$

Let us go over now to the normal atoms. Unlike positronium, they do not exist in eigenstates of C, the charge conjugation operation. The only distinguishing characteristic of the TOPE interaction remaining is T- reversal property and care must be taken to distinguish it from final state interaction effects.

The local TOPE electron-nucleon interaction in atoms has been considered in ref. [6]. Here again we will be interested mainly in the long-range TOPE interaction. Its nonrelativistic limit looks as follows:

$$V = -\frac{1}{4m_e m_p}\{(\beta_{eN} + \beta_{Ne})\frac{1}{2}(\sigma_m \sigma_n^N + \sigma_n \sigma_m^N)(p_m \frac{r_n}{r^3} + \frac{r_n}{r^3} p_m) + (\beta_{eN} - \beta_{Ne})[\sigma \times \sigma^N] \cdot \mathbf{l} \frac{1}{r^3}\} \quad (66)$$

where β_{eN} and β_{Ne} are the TOPE electron-nucleon dimensionless constants, σ and σ^N are the electron and nucleon spins, and $\mathbf{p}$ is the momentum operator.

Let us go over now to the TOPE photon-fermion interaction in an atom, confining ourselves to amplitude (13) which is of the lowest possible dimension. We will consider the case when one of the field strengths refers to the interatomic Coulomb field and so we are dealing with a TOPE one-photon emission or absorption amplitude. If the interaction discussed refers to electron, the amplitude is reduced to

$$-4A_e \frac{m_e^2}{m_p^6} i\omega Z \mu \sigma_m [\delta_{mn} \frac{4\pi}{3} \delta(\mathbf{r}) - \frac{3r_m r_n - \delta_{mn} r^2}{r^5}] B_n. \quad (67)$$

Here $\mu = |e|/2m_e$ is the Bohr magneton, Z is the nuclear charge, ω and $\mathbf{B}$ are the frequency and magnetic field of the emitted photon. We have taken into account here that the interaction is concentrated at short distances, so the Coulomb field is essentially that of an unscreened nucleus.

When dealing with the photon-nucleon interaction, the analogous emission amplitudes is

$$2A_p \frac{1}{m_p^4} i\omega \mu^N \sigma_m^N [\delta_{mn} \frac{4\pi}{3} \delta(\mathbf{r}) - \frac{3r_m r_n - \delta_{mn} r^2}{r^5}] B_n. \quad (68)$$

where $\mu^N = |e|/2m_p$ is the nuclear magneton. The explanation of the extra factor $-1/Z$ as compared to the previous formula, is that here the electric field is created by one valence electron, but not by Z protons. An overall factor 2, instead of 4, can be traced back to a somewhat different nonrelativistic reduction.

In atoms the variety of mixing states is clearly wider than in positronium. To optimize for the "best" atomic system to observe TOPE effects, high Z systems would seem to to be preferred. However, any pair of states mixed by TOPE eN interaction, also is mixed by the hyperfine interaction which exhibits the same r-dependence and Z-dependence. Since the hyperfine induced transitions serve as an irreducible background to any TOPE induced transitions, small TOPE amplitudes must be extracted by interference either with the hyperfine induced amplitude or another larger amplitude. Thus, for the long-range interaction, the Z of the atom is of little importance and other experimental factors will govern the choice of atomic system. The same conclusion holds for the TOPE photon-fermion interaction, at least for that described by formula (68).

A local TOPE eN interaction increases with Z much more rapidly than the hyperfine interaction, so in this case heavy atoms are attractive candidates [6].

Experiments to search for both long- and short-range TOPE interactions in atoms are presently being designed, and limits on the matrix elements of TOPE mixing in the range 10^2 - 10^{-2} Hz are possible [8]. The corresponding long-range limits on β_{eN} and β_{Ne} are 10^{-8} - 10^{-12}. These constants are of the same order of magnitude as β_{eq}

and β_{qe}. So, such limits would represent a significant improvement over those inferred above from the EDM measurements.

I am grateful to P. Bucksbaum and R. Conti for numerous stimulating discussions. This work was supported by NSF grant PHY-8803718. I wish also to thank Institute for Nuclear Theory at the University of Washington for the hospitality and support during the completion of writing this text.

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