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INTERNATIONAL CENTRE FOR THEORETICAL PHYSICS

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ORDER αG_μ CORRECTIONS TO THE PARITY-VIOLATING ELECTRON-QUARK POTENTIAL IN THE WEINBERG-SALAM THEORY: PARITY-VIOLATIONS IN ONE-ELECTRON ATOMS *

Bryan W. Lynn
International Centre for Theoretical Physics, Trieste, Italy.

MIRAMARE - TRIESTE

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Abstract

We have calculated the order αG_μ corrections to the order G_μ parity-violating (PV) electron-quark potential H_{PV}^{e-q} at $q^2=0$ in the standard $SU(2)_L \times U(1)$ theory using a renormalization scheme which has M_W, M_Z and M_H (Higgs' mass) as input parameters. We then use SU(3) relations to write an effective PV electron-nucleon potential H_{PV}^{e-N} in terms of the dimensionless parameters C_{1P}, C_{1N} (nuclear spin independent) C_{2P}, C_{2N} (nuclear spin dependent) and C_{3P}, C_{3N} (electron anomalous magnetic moment dependent). For $s_\theta^2 \equiv 1-M_W^2/M_Z^2 \approx .23$ and small Higgs' mass $M_H \lesssim M_W$, the corrections to C_{1P} are negligible while those to C_{1N}, C_{2N} are quite large with corrections to C_{2P} of medium importance. The M_H-dependent corrections to these four parameters are all large for large $M_H \gtrsim 1$ TeV. C_{3P}, and C_{3N} are first induced at one loop and are small.

We show that C_{2P} and C_{2N} suffer from large uncertainties due to the strong interactions which make the theoretical interpretation of experiments designed to measure these quantities less clear than was preciously thought. The other parameters are relatively free from strong interaction uncertainties. We review the diseases of the old 1934 four-fermion theory and give an overview of the renormalization of the Weinberg-Salam theory. We discuss the possibility of high-precision experiments in one-electron atoms to measure these radiative corrections and remind the reader of a previously proposed experiment in hydrogen or deuterium which will measure C_{1P} or $C_{1P}+C_{1N}$ respectively.

We have just heard the results of M. Bouchiat's[1] beautiful and careful experiment on parity-violations in atomic Cesium which, it seems, is in agreement with the predictions of the Weinberg-Salam (WS) theory.[2] The reason for all of the excitement is that an atomic physics experiment, done at very low energy and cost, has measured something fundamental to high energy physics with small experimental error; in this case the parity-violating (PV) properties of neutral currents. This is reminiscent of the experimental vindication of QED in 1947 by the observed Lamb shift.[3] The fact that the Paris experiment agrees with WS is taken as a major triumph for $SU(2)_L \times U(1)$ unification.

The question is: to what accuracy can we push these types of experiments and still get a 'theoretically clean' prediction from the fundamental gauge theory? You have heard about the difficulty and complexity of the calculation of purely atomic effects in heavy atoms and it is reasonable to assume that these calculations cannot be pushed beyond the 10% confidence level.(1,27) This, of course, brings us to one-electron atoms and, in particular, hydrogen. Since the atomic wavefunctions are known exactly in the non-relativistic limit, there is hope that it might be possible to test the higher-order corrections to PV effects. In fact, I shall show that higher-order corrections are liable, for experiments on PV effects in hydrogen, to give a large fraction of the experimental signal but that the theoretical interpretation of the results is less clear than at first glance due to strong-interaction effects. I will report here on the one-loop radiative corrections to the effective electron-quark PV potential[4] which is presumably the origin of the electron-nucleon PV potential.

The PV electron-quark (e-q) potential is defined, modulo Fourier transforms, as the S-matrix element of elastic e-q scattering at momentum transfer $q^2 = 0$. All external fermion lines are on-mass-shell ($p^2 = -m^2$). Then, the PV e–q potential may be parametrized as follows:[5]

$$H_{PV}^{e-q} = \frac{G_\mu}{\sqrt{2}} \int d^3x \Big\{ \bar{e}\gamma_\mu\gamma_5 e \, (C_{1u}\, \bar{u}\gamma_\mu u \; + \; C_{1d}\, \bar{d}\gamma_\mu d) \\ + \; \bar{e}\gamma_\mu e \, (C_{2u}\, \bar{u}\gamma_\mu\gamma_5 u \; + \; C_{2d}\, \bar{d}\gamma_\mu\gamma_5 d) \\ + \; \bar{e}\sigma_{\mu\nu} q_\nu e \, (C_{3u}\, \bar{u}\gamma_\mu\gamma_5 u \; + \; C_{3d}\, \bar{d}\gamma_\mu\gamma_5 d) \Big\} \qquad (1)$$

Here, e, u and d are *on-shell free* electron, u-quark and d-quark Dirac spinors and the Dirac algebra, as well as all other conventions[6] used, is as in Diagrammar with $\sigma_{\mu\nu} = \frac{1}{2}[\gamma_\mu, \gamma_\nu]$. $G_\mu = (1.16632 \pm .00002) \times 10^{-5}\,\mathrm{GeV}^{-2}$ is the muon decay constant. The dimensionless parameters $C_{1u}, C_{1d}, C_{2u}, C_{2d}, C_{3u}$ and C_{3d} are the object of our calculation; we want them to 1% accuracy.

H_{PV}^{e-q} has the form of a current–current interaction and could be a simple-minded generalization of Fermi's 1934 theory of weak interactions in which we introduce neutral currents as well as charged ones.[7] So what has all of the elaborate structure of gauge theories, which we learned in Prof. Hayot's lectures , bought us? To see the answer, we must look not at the form of H_{PV}^{e-q} but where it comes from.

The old four-fermion theory suffers from two incurable diseases. These are the violation of the unitarity condition at high q^2 and the fact that the theory is non-renormalizable. The unitarity condition is related to the optical theorum and says that for S-wave scattering the total cross section $\sigma < \pi |2q^2|^{-1}$. For $C_{3u} = C_{3d} = 0$ the Fermi theory gives matrix elements $\mathcal{M} \sim G_\mu$ (at tree level) independant of q^2 . However $d\sigma/d\Omega \sim |\mathcal{M}|^2 q^2$ so that for $|q^2| \sim (300\ \mathrm{GeV})^2$ the probability of interaction becomes greater than one! Because the matrix elements for massive boson exchange in a spontaneously broken gauge theory (SBGT) fall off as q^{-2} for large q^2, this is not a problem for SBGT's and the unitarity condition is obeyed.

Perhaps the most compelling reason (for theorists) to prefer a

gauge theory to Fermi's theory is renormalizability: the ability to handle systematically the ultra-violet (UV) divergences which appear in higher-order corrections to physical processes. Let's say we depict the effect of H_{PV}^{e-q} for e–u scattering by the Feynman diagram of Fig. 1a and we want to calculate the $\mathcal{O}(G_\mu{}^2)$ corrections to this process. These would be given, in part, by the diagram of Fig. 1b. This 'graph' (for very deep reasons which Prof. Stora has hinted at) is gotten by integration over products of 'time-ordered products' or 'propagators' and vertices formed according to the 'Feynman rules'.[7] For electrons, the propagator is $\overline{e\,\bar{e}} = (-i\not{p} + m_e)/(p^2 + m^2 - i\epsilon)$ and the contribution of the graph in Fig. 1b is

$$\mathcal{M}_{1b} = \int \left[\frac{-i(\not{p}_e' - \not{p}_u + \not{k}) + m_e}{(p_e' - p_u + k)^2 + m_e^2 - i\epsilon}\right]_{\alpha\beta} \left[\frac{-i\not{k} + m_u}{k^2 + m_u^2 - i\epsilon}\right]_{\delta\lambda} d^4k \; \mathcal{O}_{\alpha\beta}^{\delta\lambda}\, G_\mu^2 \qquad (2)$$

where $\mathcal{O}_{\alpha\beta}^{\delta\lambda}$ is a product of Dirac matrices and spinors depending only on external momenta and spins. Modulo all sorts of subtleties *from the theory of* complex variables this integral blows up for $\Lambda \geq k \gg$ (external momenta) as $\int^\Lambda k\,dk \sim \Lambda^2$ where Λ is a UV cut-off momentum which we would like to let become infinite. But then the matrix element would become infinite! This is the UV catastrophe.

It is the great success of 'renormalizable' field theories that these UV divergences can all be absorbed into un-measurable non-physical parameters so that the measurable physical quantities (like mass, electric charge, etc.) are all finite as $\Lambda \to \infty$. It turns out that one of the necessary conditions for this to happen is that no parameter appearing in the theory have dimensions of $(\mathrm{mass})^{-n}$ with $n > 0$. Since $G_\mu \sim (\mathrm{mass})^{-2}$ the four-fermion theory is non-renormalizable and is stripped of predictive power in higher-order corrections. Thus, radiative corrections would be an area of genuine disagreement between the (phenomenological) Fermi theory and the (renormalizable) gauge theory. An experimentally verified radiative correction would be a major triumph for non-Abelian gauge theories and would justify all of this theoretical work. After all , both the anomalous magnetic moment of the muon and the Lamb shift are radiative corrections and are two of the most compelling reasons to believe QED. An experiment capable of measuring, say, C_{1u} to 1% accuracy would test the $\mathcal{O}(\alpha)$ corrections to this quantity. Since these radiative corrections, which were the reason for the invention of gauge theories in the first place, depend on the full SBGT structure of the gauge theory (Higgs' ,gauge particles, extra fermions, etc.),this would provide the first evidence that the fundamental local gauge-symmetric structure upon which most of recent theoretical ideas are based is, in fact, related to physical experimental reality.

We will calculate the parameters $C_{1u}, C_{1d}, C_{2u}, C_{2d}, C_{3u}, C_{3d}$ through $\mathcal{O}(\alpha)$ in the 'standard' $SU(2)_L \times U(1)$ theory due to A.Salam and S. Weinberg neglecting strong interactions of the quarks. To lowest order in the fine structure constant $\alpha = (137.035963 \pm .000015)^{-1}$ the Hamiltonian H_{PV}^{e-q} is the $q^2=0$ limit of the diagram in Fig. 2 and we have $C_{1u}^T = R(1-2v_\theta)/6$, $C_{1d}^T = R(v_\theta - 2)/6, C_{2d}^T = -C_{2u}^T = Rv_\theta/2$, $C_{3u}^T = C_{3d}^T = 0$ with $v_\theta = 4s_\theta - 1, s_\theta^2 = \sin^2\theta_{WS}$ $= 1 - M_W^2/M_Z^2 = 1 - c_\theta^2$ with θ_{WS} the Z-A mixing angle and M_W and M_Z the experimental masses of the (soon to be discovered) $W^\pm$ and Z bosons and A the photon.[8] The overall factor $R = (37.281\text{GeV}/(M_W s_\theta))^2$ is expected to be very close[30] to 1 . The 'T' superscript indicates that these C^T's came from the 'tree' diagram of Fig. 2.

To the next order in α, the one-loop level, the calculation of C_{1u}^1, $C_{1d}^1, C_{2u}^1, C_{2d}^1, C_{3u}^1, C_{3d}^1$ ('1' for 'one-loop') is quite complicated due to the SBGT structure of the WS theory. When expressed in terms of the renormalized measurable quantities M_W, M_Z, m_f, M_H and α (M_H is the Higgs' scalar mass and m_f is a generic name for a fermion mass), these parameters are UV and infra-red (IR) finite and correct through $\mathcal{O}(\alpha)$. Note that the above definition of s_θ^2 is the usual one (ratio of charged to neutral current neutrino scattering) at tree level but different at the one-loop level.[9]

It is beyond the scope of this lecture to explain the details of the renormalization of $SU(2)_L \times U(1)$ but I will try to indicate the

general procedure for one-loop calculations in WS.[2,4,10]

First we make a list of all multiplets of fields appearing in the theory before spontaneous symmetry breaking (SSB). For simplicity, we'll stick to one 'generation' of fermions. All fields transform as singlets under U(1) hypercharge while under SU(2) weak isospin we have i) one left-handed (LH) lepton doublet $\ell_L^o = \begin{pmatrix} e^o \\ \nu^o \end{pmatrix}_L$ and one LH quark doublet $q_L^o = \begin{pmatrix} u^o \\ d^o \end{pmatrix}_L$ ii) four right-handed(RH) fermion singlets e_R^o , ν_R^o , u_R^o and d_R^o , iii) a triplet of SU(2) gauge vector bosons $\underset{\sim}{W}_\mu^o$, iv) a singlet U(1) gauge vector boson B_μ^o , v) a (complex) doublet of Higgs' scalars $\hat\phi^o = \frac{1}{\sqrt{2}}\begin{pmatrix} \phi_1^o + i\phi_2^o \\ H^o - i\phi_3^o \end{pmatrix}$,vi) a triplet and singlet of 'ghost' fields needed in $R_\xi(\xi=1)$ gauge to quantize the theory and subtract out extraneous degrees of freedom, which we label collectively as η_j^o, ξ_j^o (j=1,,,4).[11] The little 'o' superscript indicates that these are the unphysical bare fields while a left or right-handed fermion is gotten from the fermion field by multiplication from the left by $\gamma_L = (1+\gamma_5)/2$ or $\gamma_R = (1-\gamma_5)/2$ respectively.

Now write the 'bare' Lagrangian L^b in terms of these bare fields and some bare coupling constants g_o, g_o', λ_o and E_o^2 and bare fermion masses m_e^o, m_u^o and m_d^o according to Prof. Hayot's instructions so that the whole thing has $SU(2)_L \times U(1)$ symmetry:

$$L^b = L^b_{\psi W} + L^b_W + L^b_{\phi W} + L^b_{\psi\phi} - V^b_\phi \tag{3.1}$$

$$\begin{aligned} L^b_{\psi W} = &- \bar\ell_L^o(\partial_\mu - \tfrac{i}{2}g_o \underset{\sim}{\sigma}\cdot\underset{\sim}{W}_\mu^o + \tfrac{i}{2}g_o' B_\mu^o)\gamma_\mu \ell_L^o - \bar e_R^o(\partial_\mu + ig_o' B_\mu^o)\gamma_\mu e_R^o \\ &- \bar q_L^o(\partial_\mu - \tfrac{i}{2}g_o \underset{\sim}{\sigma}\cdot\underset{\sim}{W}_\mu^o + \tfrac{i}{6}g_o' B_\mu^o)\gamma_\mu q_L^o - \bar d_R(\partial_\mu + \tfrac{i}{3}g_o' B_\mu^o)\gamma_\mu d_R^o \\ &- \bar u_R^o(\partial_\mu - \tfrac{2}{3}ig_o' B_\mu^o)\gamma_\mu u_R^o \end{aligned} \tag{3.2}$$

$$L^b_W = -\tfrac{1}{4}(\partial_\mu B_\nu^o - \partial_\mu B_\nu^o)^2 - \tfrac{1}{4}(\partial_\mu \underset{\sim}{W}_\nu^o - \partial_\nu \underset{\sim}{W}_\mu^o + g_o \underset{\sim}{W}_\mu^o \times \underset{\sim}{W}_\nu^o)^2 \tag{3.3}$$

$$L^b_{\phi W} = -\left|(\partial_\mu - \tfrac{i}{2}g_o \underset{\sim}{\sigma}\cdot\underset{\sim}{W}_\mu^o - \tfrac{i}{2}g_o' B_\mu^o)\hat\phi^o\right|^2 \tag{3.4}$$

$$\begin{aligned} L^b_{\psi\phi} = &- m_e^o E^{-1}\{\bar\ell_L^o \phi^o e_R^o + \bar e_R^o \phi^{o\dagger}\ell_L^o\} - m_u^o E^{-1}\{\bar q_L^o \hat\phi^o u_R^o + \bar u_R^o \hat\phi^{o\dagger} q_L^o\} \\ &- m_d^o E^{-1}\{\bar q_L^o \phi^o d_R^o + \bar d_R^o \phi^{o\dagger} q_L^o\} \end{aligned} \tag{3.5}$$

$$V^b_\phi = \lambda_o(\phi^o \phi^{o\dagger} - E_o^2)^2/2 \tag{3.6}$$

with $\underset{\sim}{\sigma}$ the usual Pauli matrices and $\hat{\phi}^o = -i\sigma_2\phi^{o*}$. The total Lagrangian L_{tot} is gotten by adding the gauge-fixing and 'ghost' Lagrangians $L_{g.f.}$ and L_{ghost} in such a way that L_{tot} has the so-called BRS symmetry.(11)

It is the fundamental result of renormalization theory that all UV divergences appearing in a given calculation in a renormalizable field theory are fictitious and due to the fact that there is no absolute momentum (or any other) scale in the universe but only ratios of momenta to some arbitrary standard. Therefore, when we rescale the bare fields

$$\begin{array}{lll} \underset{\sim}{W}^o_\mu = Z_W^{1/2}\,\underset{\sim}{W}_\mu & \ell^o_L = (Z^\ell_L)^{1/2}\,\ell_L & q^o_L = (Z^q_L)^{1/2}\,q_L \\ B^o_\mu = Z_B^{1/2}\,B_\mu & \nu^o_R = (Z^\nu_R)^{1/2}\,\nu_R & u^o_R = (Z^u_R)^{1/2}\,u_R \\ \phi^o = Z_\phi^{1/2}\,\phi & e^o_R = (Z^e_R)^{1/2}\,e_R & d^o_R = (Z^d_R)^{1/2}\,d_R \end{array} \qquad (4)$$

and the bare coupling constants and masses

$$\begin{array}{ll} (Z^\ell_L)^{1/2}(Z^e_R)^{1/2}Z_\phi^{1/2}m^o_e = m_e+\delta m_e & Z_W^{1/2}g_o = g+\delta g \\ (Z^q_L)^{1/2}(Z^u_R)^{1/2}Z_\phi^{1/2}m^o_u = m_u+\delta m_u & Z_B^{1/2}g'_o = g' \\ (Z^q_L)^{1/2}(Z^d_R)^{1/2}Z_\phi^{1/2}m^o_d = m_d+\delta m_d & Z_\phi^{-1}E^2_o = E^2+\delta E^2 \\ & Z^2_\phi\lambda_o = \lambda+\delta\lambda \end{array} \qquad (5)$$

all UV divergences occuring in matrix elements or Green's functions may be absorbed into the Z's and the δ's (wavefunction renormalization and coupling constant renormalization respectively). Note that $\delta g' = 0$ because of the U(1) hypercharge symmetry.(12) In general, we would also have to rescale the ghost fields but the first effects of this would only appear at the two-loop level $\mathcal{O}(\alpha^2)$ and we will disregard ghost field renormalization for the one-loop calculation of interest here. If we write $Z_j = 1+\delta Z_j$ for each of the Z's we may divide the quantities appearing in L^b into three sets: $\{\delta\} = \{\delta Z^\ell_L, \delta Z^q_L, \delta Z^e_R, \delta Z^u_R, \delta Z^d_R, \delta Z_W, \delta Z_B, \delta Z_\phi, \delta g, \delta\lambda, \delta E^2, \delta m_e, \delta m_u, \delta m_d\}$, the renormalized fields $\{\phi\} = \{e, \nu, u, d, \underset{\sim}{W}_\mu, B_\mu, \phi\}$ and the renormalized parameters $g_i = \{g, g', \lambda, E^2, m_e, m_u, m_d\}$. $\{\delta\}$ is $\mathcal{O}(\alpha)$ and formally infinite (but contains finite parts as well which are very important to maintain self-consistency) while $\{\phi\}$ and g_i are finite, $\mathcal{O}(1)$ and we have dropped only $\mathcal{O}(\alpha^2)$ terms in their definitions.

So far, the seven constants g_i are just parameters in a long mathematical expression and have no physical meaning. In order to give them physical content, we appeal to an experimental fact; there are <u>no absolute experimental measurements</u>, <u>only ratios of experimental results</u>. This is because there are no absolute scales (or units) of energy, distance, electric charge or anything else. Some 'standard' experiments are done to define 'standard' scales (and units) and the results of all other experiments are expressed in terms of these standards. Thus, we could <u>define</u> the electron-volt so that $m_e = .5110034$ MeV and compare all subsequent measurements with this standard.

The need for renormalization is not a fiction or even a theoretical drawback but a reflection of the philosophical unavailability of absolute fundamental units. Thus, the absolute result of the calculation of an S-matrix element in terms of g_i and $\{\delta\}$ by itself makes no physical sense. It is only ratios of S-matrix elements which have physical content because, in these, the arbitrary rescaling in eq. (4) and (5) will drop out of our prediction for the ratios of physical experimental measurements. Actually, because the units of the parameters in g_i and $\{\delta\}$ are not all the same, we do not form ratios of matrix elements but instead <u>we choose some 'standard' matrix elements and write all the others in terms of these</u>. <u>This renders the entire S-matrix UV finite</u>, independent of $\{\delta\}$ and only implicitly dependent on g_i through the seven 'standard' matrix elements.

The procedure is then the following; Calculate 7 standard matrix elements $S_j(g_i, \{\delta\}, \Lambda)$ where Λ is the UV cut-off momentum and $i,j=1,\ldots,7$. (Λ enters here because, as we will see, S_j will be defined by sets of 'Feynman graphs' generated according to the 'Feynman rules' of the theory and these will involve integrals over internal loop-momenta, as in eq. (2), which will be ultra-violet (UV) divergent.)(13)

Next, invert these equations to write the $g_i(S_j, \{\delta\}, \Lambda)$. Now calculate any other S-matrix element, say, $S_8(g_i, \{\delta\}, \Lambda)$ in terms of g_i, $\{\delta\}$ and Λ and then solve for S_8 in terms of S_j, $\{\delta\}$ and Λ by eliminating the g_i. The theory is constructed so that, after SSB, the functional dependence on $\{\delta\}$ and Λ, formally infinite quantities, will cancel! Thus, we may write S_8, or **any** other S-matrix element, in terms of S_j ($j=1,,,7$) alone and there appear no more troublesome UV infinities. A slightly more complicated procedure can be made to render all Green's functions (external momenta off-mass-shell) UV finite but there we must give some prescription for the $\{\delta\}$ as well.[4]

Let's see how this works in practice. We start with the bare Lagrangian L^b and divide it into 3 parts for convenience. Two parts are UV finite and one is (formally) infinite; $L^b = L^r_{quadratic} + L^r_{vertices} + L_{c.t.}$ ('r' for 'renormalized' and 'c.t.' for 'counter-terms'). Here, the renormalized Lagrangian L^r is gotten from L^b in eqs. (3.1) to (3.6) by dropping all of the little 'o' superscripts or subscripts or, equivalently, setting all the $\{\delta\} = 0$. We include in $L^r_{quadratic}$ all those terms in L^r which are products of exactly 2 fields (fermion, vector gauge boson or Higgs' scalar); it will give the propagators. Everything else in L^r is thrown into $L^r_{vertices}$ and this gives renormalized interaction vertices used to draw our Feynman graphs. $L_{c.t.}$ will give another set of interaction vertices called the 'counter-term insertions' which will exactly subtract out all UV infinities generated by integrals over internal loop momenta in the Feynman graphs constructed using L^r.

An example will show the simplicity of this procedure. Look at the lepton-gauge boson part of $L^b_{\psi W}$ in eq. (3.2) and expand it out in $\{\delta\}$, consistently dropping terms of $\mathcal{O}(\delta^2) \sim \mathcal{O}(\alpha^2)$ and all quark terms:

$$\begin{aligned} L^b_{\psi W} = & - \bar{e}_R(\partial_\mu \gamma_\mu) e_R - \bar{\ell}_L(\partial_\mu \gamma_\mu)\ell_L + \tfrac{i}{2}\bar{\ell}_L(g\underset{\sim}{\tau}\cdot \underset{\sim}{W}_\mu - g' B_\mu)\gamma_\mu \ell_L \\ & - ig'\,\delta Z^e_R \bar{e}_R B_\mu \gamma_\mu e_R - \tfrac{i}{2}\bar{\ell}_L\{-2i\delta Z^\ell_L \partial_\mu - (\delta g + g\delta Z^\ell_L)\underset{\sim}{\tau}\cdot\underset{\sim}{W}_\mu + g'\delta Z^\ell_L B_\mu\}\gamma_\mu \ell_L \\ & + \mathcal{O}(\alpha^2) + \text{quark terms} \end{aligned} \qquad (6)$$

The 1st and 2nd terms go into $L^r_{\psi W;quadratic}$ and, together with the lepton-mass terms from $L^r_{\psi\phi;quadratic}$, give the renormalized propagators for the ν ($m_\nu=0$) and the electron. The 3rd term gives the renormalized lepton-gauge boson vertices and goes into $L^r_{\psi W;vertices}$ while the 4th and 5th terms go into $L_{\psi W;c.t.}$ and give the lepton self-energy insertions and lepton-gauge boson vertex insertions which will subtract out all the UV divergences occuring, to one loop, in the effective lepton self-energy and lepton-gauge boson couplings.

Finally, we induce SSB by giving the renormalized scalar Higgs' field a vacuum-expectation-value $\langle 0|H|0\rangle = \sqrt{2}E$ and introduce the WS angle $s_\theta = g'(g^2+g'^2)^{-1/2}$ and the electric charge $e = s_\theta g = c_\theta g'$ and rotate the fields B_μ and W^3_μ to give the $Z_\mu = c_\theta W^3 - s_\theta B_\mu$ and the photon $A_\mu = s_\theta W^3_\mu + c_\theta B_\mu$. The result of this very long but systematic procedure is a list of the renormalized propagators and vertices and the counter-term insertions called the 'Feynman rules' of the theory. These are shown pictorially for the lepton-gauge boson interactions coming from $L^b_{\psi W}$ in eq. (6) in Fig. 3. All of the possible topologies occuring in the Feynman rules resulting from the expansion of the complete bare Lagrangian (including the ghosts) are shown in Fig 4. In Figs. 3 and 4, the solid lines with arrows indicate fermion propagators $\psi = \nu_e, e, u, d$, etc. while the wiggly lines are gauge bosons $W_\mu = W^\pm_\mu, Z_\mu, A_\mu$, dashed lines are Higgs' scalars $\phi = \phi^\pm, \phi_3, H$ and dotted lines with arrows are ghosts χ_j, η_j ($j=1,,,4$). The renormalized propagators and vertices have no crosses while the counter-term insertions are indicated by crosses. There are 14 propagators, 69 vertices and who knows how many counter-term insertions in $SU(2)_L \times U(1)$.[14] By comparison, QED has 2 propagators, 1 vertex and 3 counter-term insertions!

We are now in a position to calculate our 7 'standard' matrix elements in terms of which we will express the rest of the S-matrix. It turns out that it is possible to form combinations of the g_i so that these combinations themselves are the results of predictions for experiments. They are

$d_i = m_e,\ m_u,\ m_d,\ M_W = gE/\sqrt{2},\ M_Z = E(g^2 + g'^2)^{1/2}/\sqrt{2},\ M_H^2 = 2\lambda E^2,$

and $\alpha = e^2/(4\pi) = g^2 g'^2/(4\pi(g^2+g'^2))$. We need only to write the entire S-matrix in terms of these and it is guarranteed to be UV finite, gauge invariant and independent of Λ and $\{\delta\}$. We take the numerical values of these quantities as experimental input data ('d_i' for 'data') in our calculation of the PV electron-quark potential and will express the C's in terms of the d_i. The choice of which data or matrix elements to use as our 'standard' ones is called the 'renormalization prescription' and it must be emphasized that this particular choice is quite arbitrary; it just depends on which data from previous experiments in terms of which one chooses to predict the numerical values of $C_{1u}, C_{1d}, C_{2u}, C_{2d}, C_{3u}$ and C_{3d}. We choose this prescription because we believe that, upon discovery of the $W^{\pm}$ and Z, the above will be the most accurately know data around and also the most convenient for direct comparison of the results of low energy atomic physics experiments, high energy leptonic(and other) experiments and the precise values of the muon lifetime and anomalous magnetic moment.[8,9]

We now form the one-particle-irreducible (1PI) vertex and self-energy parts which are the backbone of the theory. These are indicated by the shaded blobs in Fig. 5 and are gotten by putting together all possible combinations of the renormalized propagators and vertices consistent with the Feynman rules in such a way that the diagram cannot be cut into 2 distinct parts by cutting one particle line, adding the appropriate counter-term insertion(s) and then integrating over internal loop momenta k, as in eq. (2), with momentum conservation at the vertices. The Feynman diagrams with an internal loop will, in general, have a UV divergence $\sim \ln(\Lambda)$ as the UV cut-off momentum $\Lambda \gg k; \Lambda \to \infty$ but this will be exactly cancelled by the $\ln(\Lambda)$ divergences appearing in the definitions of $\{\delta\}$ in the counter-term insertion graphs after we have defined our renormalization prescription. Then, all 1PI vertex and self-energy parts will be UV finite . The results will be unique because both the divergent and finite parts were generated in a systematic way from our starting point L^b! The reason that theorists are so excited about SBGT's is that

<u>one can eliminate all UV divergences in all orders in α by the rescaling in eqs. (4) and (5) when one rescales the ghost fields as well!</u>[15] This is, to say the least, a nifty trick.

One of the most difficult things about this type of calculation is to check that the results are correct. Given the length and complexity of one-loop calculations in $SU(2)_L \times U(1)$, this is no mean task. Since all of the 1PI parts(and,indeed, all S-matrix elements and Green's functions) come from L^b in eq.(3.1) one can imagine that there are many relations between the results of <u>different</u> calculations. For example, the 1PI Z-A mixing is related to the 1PI ϕ_3-A mixing because the ϕ_3 is, in some sense, the longitudinal component of the massive vector Z boson. These relations are called the Ward-Takahashi-Slavnov-Taylor[10] identities; they are true even non-perturbatively and follow from the fundamental $SU(2)_L \times U(1)$ symmetry of the theory. We may use them as a very strong check on the correctness of our results simply by calculating the relevant 1PI parts separately and making sure that the known relations between them (the Ward identities) are satisfied. Further, the fact that all quadratic divergences $\sim \Lambda^2$ cancel even before renormalization (addition of the counter-term graphs) and that all logarithmic divergences cancel after renormalization (even in Green's functions), along with other checks not mentioned here, gives us added confidence in the correctness of our formulae for the C's. We emphasize that, since it is almost impossible to check this type of calculation without doing it yourself, the substantiation of one's claim that the results are correct is every bit as important as the result itself.

Having formed the UV finite 1PI self-energy and vertex parts, we can now calculate the C^i's in H_{PV}^{e-q}. We will calculate H_{PV}^{e-q} only through $\mathcal{O}(\alpha G_\mu)$ discarding terms of $\mathcal{O}(\alpha^2 G_\mu)$, $\mathcal{O}(G_\mu^2)$, $\mathcal{O}(\alpha G_\mu m_e/m_h)$, $\mathcal{O}(G_\mu m_h^2/M_H^2)$ where m_h is some typical hadronic mass. We take the limit $q^2=0$ and assume $M_H \gtrsim 3$GeV or so.[16] Further, we neglect all Cabibbo-type mixings between generations of quarks[17] and neglect completely the effects of strong interactions of the quarks.

The C^1's come from the contributions of the one-loop graphs, displayed in terms of the 1PI parts, of Figs. 6,7 and 8. We find, after adding the contributions from the tree graph in Fig. 2 (the C^T's), the parameters appearing in H_{PV}^{e-q} in eq. (1) correct through $\mathcal{O}(\alpha)$ to be [4]

$$C_{1u} = \frac{1-2v_\theta}{6} R + \frac{\alpha}{6\pi}\Big\{14 - \frac{9v_\theta}{2} - \frac{25}{6s_\theta^2} - \frac{25v_\theta}{16s_\theta^2c_\theta^2} - 3v_\theta \ln\frac{M_Z^2}{m_h^2} + \frac{2v_\theta}{3}\ln\frac{M_Z^2}{m_e^2} - \Big(\frac{5}{3} + \frac{c_\theta^4}{s_\theta^4}\Big) N \ln c_\theta^2 + \frac{1}{3}\sum_{i=1}^{N} \ln\frac{m_{d_i}^2}{m_{\ell_i^+}^2} + \Pi_1 + \Pi_2 \Big\} \tag{7.1}$$

$$C_{1d} = \frac{v_\theta - 2}{6} R + \frac{\alpha}{6\pi}\Big\{-\frac{61}{6} + \frac{3v_\theta}{4} + \frac{17}{3s_\theta^2} - \frac{5v_\theta}{16s_\theta^2c_\theta^2} - \frac{3v_\theta}{2}\ln\frac{M_Z^2}{m_h^2} - \frac{v_\theta}{3}\ln\frac{M_Z^2}{m_e^2} + \frac{20N}{9} + \Big(\frac{5}{3} + \frac{c_\theta^4}{s_\theta^4}\Big) N \ln c_\theta^2 - \frac{2}{3}\sum_{i=1}^{N}\ln\frac{M_W^2}{m_{\ell_i^+}^2} - \frac{2}{3}\sum_{i=1}^{N}\ln\frac{M_W^2}{m_{u_i}^2} - \Pi_1 - 2\Pi_2\Big\} \tag{7.2}$$

$$C_{2u} = -\frac{v_\theta}{2} R + \frac{\alpha}{6\pi}\Big\{\frac{169}{27} - \frac{275v_\theta}{54} - \frac{11}{3s_\theta^2} - \frac{17v_\theta}{16s_\theta^2c_\theta^2} + \Big(\frac{13}{9} - \frac{14v_\theta}{9}\Big)\ln\frac{M_Z^2}{m_h^2} + \frac{2}{3}\ln c_\theta^2 + \frac{20N}{9} - \Big(\frac{5}{3} + \frac{c_\theta^4}{s_\theta^4}\Big) N \ln c_\theta^2 - \frac{2}{3}\sum_{\substack{\text{light}\\ \text{quarks}}}^{2N} \ln\frac{M_W^2}{m_{q_i}^2} + \Pi_1\Big\} \tag{7.3}$$

$$C_{2d} = \frac{v_\theta}{2} R + \frac{\alpha}{6\pi}\Big\{-\frac{248}{27} + \frac{655v_\theta}{108} + \frac{17}{3s_\theta^2} - \frac{5v_\theta}{16s_\theta^2c_\theta^2} - \Big(\frac{1}{9} + \frac{11v_\theta}{18}\Big)\ln\frac{M_Z^2}{m_h^2} - \frac{4}{3}\ln c_\theta^2 - \frac{20N}{9} + \Big(\frac{5}{3} + \frac{c_\theta^4}{s_\theta^4}\Big) N \ln c_\theta^2 + \frac{2}{3}\sum_{\substack{\text{light}\\ \text{quarks}}}^{2N} \ln\frac{M_W^2}{m_{q_i}^2} - \Pi_1\Big\} \tag{7.4}$$

$$C_{3u} = -C_{3d} = \frac{\alpha v_\theta}{8\pi} \tag{7.5}$$

with $N=3$ the number of light fermion generations, $m_h = m_u$ in the absence of strong interactions, $m_{\ell_i^+}$, m_{u_i}, m_{d_i} the masses respectively of the i^{th} charged lepton, up-type quark and down-type quark. Π_1 and Π_2 come from combinations of Π_{ZA} (from the Z-A mixing graphs of Fig. 7) and Π_{ZZ} (from the Z-Z self-energy graphs of Fig. 6) and depend in an algebraically complicated way on $s_\theta^2 \equiv 1 - M_W^2/M_Z^2$ and $\rho = M_H^2/M_W^2$.

$$\Pi_2 = \Pi_{ZZ} - \Pi_{ZA}/3 - 74s_\theta^2/9 \qquad \Pi_{ZA} = \Pi_{ZA}^W + \Pi_{ZA}^H$$
$$\Pi_1 = (\Pi_{ZA} - 3v_\theta\Pi_2)/(4s_\theta^2) . \qquad \Pi_2 = \Pi_2^W + \Pi_2^H \tag{8.1}$$

$$\Delta(p^2, m_1, m_2) = \Big[B_2 + B_1/2 + B_0 + \frac{\rho}{2}(B_1 + B_0/2)\Big](p^2, m_1, m_2) + \frac{13}{12}\ln M_W^2 + \frac{\rho}{8} \tag{8.2}$$

$$\Pi_2^H = \mathrm{Re}\Big\{-\Delta(-M_W^2, M_H, M_W)\Big\} + \frac{1}{c_\theta^2}\ln c_\theta^2 + \Big(7 + \frac{6\rho c_\theta^2}{1-\rho c_\theta^2}\ln\rho c_\theta^2\Big)/(8c_\theta^2) \tag{8.3}$$

$$\Pi_{ZA}^H = \mathrm{Re}\Big\{\frac{3c_\theta^2}{s_\theta^2}\Delta(-M_W^2, M_H, M_W) - \frac{3}{s_\theta^2}\Delta(-M_Z^2, M_H, M_Z) + \frac{3\rho}{8} + \frac{3\rho}{2}(B_1 + B_0/2)(-M_Z^2, M_H, M_Z)\Big\} \tag{8.4}$$

$$\Pi_2^W = -\mathrm{Re}\Big\{\Big[(8c_\theta^2+1)\Big(B_2 + \frac{1+c_\theta^2}{2c_\theta^2}B_1\Big) + \Big(\frac{5}{4c_\theta^2} - 4c_\theta^2 - 1\Big)B_0\Big](-M_W^2, M_Z, M_W)\Big\} + \Big(\frac{10c_\theta^2}{3} + \frac{35}{12} - \frac{1}{c_\theta^2}\Big)\ln M_W^2 \tag{8.5}$$

$$\Pi_{ZA}^W = -\frac{3c_\theta^2}{s_\theta^2}\Pi_2^W - \frac{3c_\theta^2}{s_\theta^2}\mathrm{Re}\Big\{\Big[\Big(8c_\theta^2 + \frac{(1-2s_\theta^2)^2}{c_\theta^2}\Big)(B_2 + B_1/2) - 2c_\theta^2 B_1 + (2-9c_\theta^2)B_0\Big](-M_Z^2, M_W, M_W)\Big\} + \Big(-37 + 21s_\theta^2 + \frac{63}{4s_\theta^2}\Big)\ln M_W^2 \tag{8.6}$$

$$B_n(p^2, m_1, m_2) = \frac{(-1)^{n+1}}{n+1}\Big\{\ln(-p^2 - i\epsilon) + \sum_{j=1}^{2}\Big[\ln(1-x_j) + F(n+1, x_j)\Big]\Big\} \tag{8.7}$$

where x_j are the roots of the equation $-p^2 x^2 + (p^2 + m_2^2 - m_1^2)x + m_1^2 = 0$ and

$$F(n,x) = -x^n \ln\Big(\frac{x-1}{x}\Big) - \sum_{k=1}^{n}\frac{x^{n-k}}{k} \tag{8.8}$$

Some comments on these formulae are in order here. Note that we sum only over light fermions. In principle, an ultra-heavy fermion doublet $f = \binom{f_u}{f_d}$ could contribute via the 9^{th}, 10^{th} and 11^{th} graphs in the Z-A mixing in Fig. 5b (or in the Z-Z self-energy) and we could count the number of fermion generations in existence by doing an experiment to measure, say, C_{2u}. The WS theory gives the sensible answer that if the masses of the fermions in the doublet are nearly degenerate $m_{f_u} \simeq m_{f_d} \gg M_W$ (SU(2) symmetry not very badly broken by the mass matrix), the contribution of such a doublet is exactly cancelled by the counter-terms; the last diagram in Fig. 5b (and a similar counter-term diagram for the Z self-energy). Thus, the renormalization program works in such a way that a doublet of quarks weighing 1 ton is prevented from affecting low-energy physics. Note, further, that there is no singularity as the neutrino mass $m_\nu \to 0$ and that we have summed over 3 quark colors for each quark flavor in the Z-A and Z-Z 1PI parts.

Another observation is that the result of the one-loop calculation can be written as a current-current interaction as in eq. (1) in the first place. Although many such terms appeared in the intermediate stages of the calculation, all UV and IR divergences have cancelled, all terms not suppressed by an overall factor of αG_μ have cancelled and all of the long-ranged PV interactions carried by photons in Fig 7 have cancelled as well. That these 'coincidences' must occur can be proved using dispersion relations and Ward identities,[18] so that the WS theory predicts that PV interactions are weak and short-ranged, even at the one-loop level. Further, one could have imagined that scalar-pseudo scalar terms like $\bar{e}\gamma_5 e\bar{u}u$ or $\bar{e}e\bar{u}\gamma_5 u$ could have been induced at the one-loop level by the 'box diagrams' of Fig. 8. These are all suppressed by factors of m_e/m_{quark} and discarded.

Note the dependence on the mass M_H of the physical Higgs' scalar H in Π_1 and Π_2, a large fraction of which is due to the Z-A mixing counter-ter graph's (last diagram in Fig. 5b) contribution Π^H_{ZA}. Although none of the other Z-A mixing graphs have an H scalar in them, the finite part of the counter-term graph does depend on M_H because the W^+ and Z self-energies, used to calculate the input parameters M_W and M_Z, do have internal H Higgs' lines. This is an example of the complicated internal consistency requirements of a correct renormalization prescription generated by the long-winded, but straightforward, expansion of L^b as discussed before.

Finally, there were many large contribution to the C^1's but they tend to cancel rather than enhance eachother. Since the number of graphs in some very high order in α is huge (much larger than 137), this cancellation is crucial for the justification of perturbation theory at all, no matter how small the coupling constant. Thus, we should not be suprised that partial calculations, based on the examination only of 'large logs' $\sim \ln \frac{M_Z^2}{m_e^2}$

or the box diagrams of Fig. 8 give answers so far from the correct ones.[19] As you know, many small numbers can conspire to cancel a large one and this has happened repeatedly in this complete calculation of the radiative corrections to the PV electron-quark potential.

Now comes the great leap of faith. We would like to relate the C_{iu} and C_{id} $(i=1,2,3)$ for quarks to similar parameters C_{iP} and C_{iN} for protons (P) and neutrons (N) occuring in the effective electron-nucleon (e-N) Hamiltonian H^{e-N}_{PV}:

$$H^{e-N}_{PV} = \frac{G_\mu}{\sqrt{2}} \int d^3x \Big\{ \bar{e}\gamma_\mu\gamma_5 e\,(C_{1P}\bar{P}\gamma_\mu P + C_{1N}\bar{N}\gamma_\mu N) + \bar{e}\gamma_\mu e\,(C_{2P}\bar{P}\gamma_\mu\gamma_5 P + C_{2N}\bar{N}\gamma_\mu\gamma_5 N) + \bar{e}\sigma_{\mu\nu} q_\nu e\,(C_{3P}\bar{P}\gamma_\mu\gamma_5 P + C_{3N}\bar{N}\gamma_\mu\gamma_5 N) \Big\} \qquad (9)$$

For this, we need some model of the effects of strong interactions of quarks inside nucleons.

It is known that sets of hadrons can be grouped together according to SU(3) multiplets. The members of the multiplet, in this case the P, N, Σ^+, Σ^-, Σ^0, Ξ^-, Ξ^0 and Λ in an octet, all have similar masses and decay rates and magnetic moments, etc. related by SU(3) formulae to reasonably good accuracy. Presumably, this is due to a (softly broken) SU(3) global symmetry among u, d and s (strange) quarks, described by Cabibbo, which also gives the relations:[20,25]

$$C_{1P} = 2C_{1u} + C_{1d} \qquad (10.1)$$

$$C_{1N} = C_{1u} + 2C_{1d} \qquad (10.2)$$

$$C_{2P} = g_A/g_V(2C_{2u} + C_{2d}) - 2D(C_{2u} + C_{2d}) \qquad (10.3)$$

$$C_{2N} = g_A/g_V(C_{2u} + 2C_{2d}) - 2D(C_{2u} + C_{2d}) \qquad (10.4)$$

$$C_{3P} = g_A/g_V(2C_{3u} + C_{3d}) - 2D(C_{3u} + C_{3d}) \qquad (10.5)$$

$$C_{3N} = g_A/g_V(C_{3u} + 2C_{3d}) - 2D(C_{3u} + C_{3d}) \qquad (10.6)$$

with $g_A/g_V = 1.255 \pm .006$ and $D = .826 \pm .007$. These relations reflect the hypothesis that the strong interactions affect the axial-vector coupling, but not the vector coupling, in β decay: $N \to P + e^- + \bar{\nu}_e$.

Let us assume that the above, along with the substitution $m_h = m_{proton}$ in eqs. (7.1) to (7.4), properly describes the relationship between the quarks in H_{PV}^{e-q} and the nucleons in H_{PV}^{e-N}. We then have $C_{3P} = -C_{3N} = (1.255 \pm .006) R v_\theta / (8\pi)$. We list the values of C_{1P}, C_{1N}, C_{2P}, and C_{2N} for various values of $s_\theta^2 = 1 - M_W^2/M_Z^2$ and M_H in Table I. We have used[21] $M_W = 83.1$ GeV and quark masses $m_u = m_d = .4$ GeV, $m_s = .7$GeV, $m_c = 1.8$ GeV, $m_t = 5$ GeV, $m_b = 10$ GeV (?) and lepton masses $m_e = .511$ MeV, $m_\mu = 105$ MeV and $m_\tau = 1.78$ GeV.

Aside from going through a detailed inspection of Table I, we may make the following observations. At tree level,[30] the quantities $C_{1P}^T = -R v_\theta/2$, $C_{1N}^T = -R/2$, $-C_{2P}^T = C_{2N}^T = (1.255 \pm .006) R v_\theta/2$ From the ratio of the neutral current to charged current neutrino scatterings, we know that, <u>at tree level</u>,[22] $s_\theta^2 = .23 \pm .01$ so that we expect $v_\theta = 4s_\theta^2 - 1 \approx -.08$ at tree level. This means that there is an unfortunate coincidental (?) suppression of C_{1P}^T, C_{2P}^T and C_{2N}^T by an extra order of magnitude. Since radiative corrections in $SU(2)_L \times U(1)$ usually give 3% to 4% corrections to quantities which have no such coincidental suppression,[23] we expect that the one-loop corrections might account for 30% to 40% of the values of C_{1P}, C_{2P} and C_{2N} and 3% to 4% of the value of C_{1N}.

We list in Table II the ratios of the one- loop corrections to the tree level values for these parameters for $s_\theta^2 = .231$ and various values of the Higgs' mass M_H. We see that, for small Higgs' mass ($M_H = 3$ GeV), the corrections to C_{1P}^T and C_{2P}^T are 2.94% and -11.8% respectively; suprisingly small considering the suppression of these parameters by v_θ at tree level. On the other hand, C_{1N}^T and C_{2N}^T are corrected by -2.65% and -20.2% respectively which are quite large given that C_{1N}^T is <u>not</u> suppressed by v_θ. Thus, we conclude that, for small M_H, the radiative corrections are large for neutrons, small for C_{1P} and of medium importance for C_{2P}.

For very large Higgs' mass ($M_H = 10$ TeV) there are large radiative corrections to all the C's. The Higgs'-mass-dependent corrections above account for 2.6% corrections to C_{1N}^T and about 28% corrections to each of

C_{1P}, C_{2P}, and C_{2N}. For $s_\theta^2 = .231$, the largest M_H-dependent contribution is, by far, from the counter-term insertion in the Z-A mixing. Therefore, for large Higgs' mass and s_θ^2 near .23, a large part of the experimental signal in a measurement of C_{1P}, C_{2P} or C_{2N} (and, to a lesser extent, C_{1N}) is liable to come from the deepest part of the gauge-structure of the WS theory, the counter-term sector, and, in that, from the Higgs' sector which is central to the idea of spontaneous symmetry breaking (SSB). It turns out that we are not really allowed to do a perturbation expansion in the parameter λ in eq. (3.6) for $M_H \gtrsim 10\, M_W$ so that these M_H-dependent corrections are an indication of the order of magnitude rather than a proper result for $M_H \gtrsim .8$ TeV. Nevertheless, these calculations show that it is important to consider what <u>does</u> happen, for such a large Higgs' mass, to our PV parameters.[24] Who knows? An atomic physics experiment might discover the first effects of the elusive Higgs' scalar H if its' mass is too large for it to be produced in high-energy accellerators!

We must now face the most important question. Do the above results properly give the predictions of the WS theory for experiments on PV in one-electron atoms?

The largest source of uncertainty is, of course, the possible effects of strong interactions. The errors quoted for g_A/g_V and D in eqs. (10.1) to (10.6) are <u>experimental</u> errors, but one should also question the assumptions which went into these formulae and the data analysis in the first place.

The formulae for C_{1P} and C_{1N} and the terms proportional to g_A/g_V can be derived by invoking only strong SU(2) global symmetry between protons and neutrons. This is the best approximate symmetry of hadrons in nature[25] (good to better than 1% accuracy) so we may believe the quoted error in g_A/g_V (the ratio of axial-vector to vector couplings in neutron β decay). Since $C_{3u} + C_{3d} = C_{2u}^T + C_{2d}^T = 0$, we may expect that the results for C_{1P}, C_{1N}, C_{3P}, C_{3N} and the <u>tree</u> level values C_{2P}^T, C_{2N}^T are reliable.

The radiative corrections C_{2P}^1 and C_{2N}^1 suffer from at least three sources of strong-interaction uncertainty. Firstly, the value of D quoted does not reveal the fact that the chiral $SU(3) \times SU(3)$ symmetry on which it is based is a good approximate symmetry of hadrons only to 10% to 30%,[25] depending on the process considered. Thus, the error quoted in D is misleading; it could really be off by as much as 30%. Secondly, there could be purely nuclear PV effects (exchange of W^{+} and Z between quarks within the nucleus) which are carried out to the electrons by photons. This would contribute only to C_{2P} and C_{2N}. Thirdly, Wilcek et. al. have shown that the so-called axial-vector (γ_5) anomaly could change $C_{2u}, C_{2d}, C_{3u}, C_{3d}$ by 10% ($\Delta C_{2u} = \Delta C_{2d} \simeq .1\ C_{2u}$ and $\Delta C_{3u} = \Delta C_{3d} \simeq .1 C_{3u}$).[26] Since C_{3u}, C_{3d} and the tree-level values C^T_{2u}, C^T_{2d} are accidentally suppressed by a factor of v_θ, our confidence in C_{3P}, C_{3N} at the one-loop level and C^T_{2P}, C^T_{2N} at the tree level is not too badly shaken. However, given the above three (and probably more) sources of strong-interaction uncertainty, our confidence in the precise prediction of the WS theory for C_{2P}, C_{2N} to, say, 10% is ruined. Unfortunately, all of the experiments on PV in hydrogen now in progress[1] are trying to measure C_{2P} so that the theoretical interpretation of the experimental signal is liable to be quite a lot less clear at the 10% level than originally thought, due to the accidental suppression of C_{2P} by $v_\theta \simeq -.08$ at the tree level and strong-interaction effects.

There is another source of theoretical uncertainty for experiments on the PV effects in one-electron atoms which has nothing to do with the strong interactions. This is due to the fact that the electron is in a bound-state with the nucleus. We should properly use QED bound-state wavefunctions (given, to first approximation by the well-known hydrogen-like wavefunctions) as basis states around which to do perturbation theory rather than the free-particle states defining the S-matrix. We would then have to take into account an infinite number of Coulomb-photon exchanges between the quarks and the electron in Figs. 6,7 and 8 in this bound-state field theory. To make matters even worse, it is not an energy-level shift which is to be measured but a transition rate. This means that we should properly consider the $\mathcal{O}(\alpha^2 G_F)$ Lamb shift in hydrogen and take its imaginary part in order to make the connection between the radiative corrections of the WS theory and the PV effects in the transition $2S \rightarrow 1S + \gamma$ in hydrogenic atoms. The reason is that, for bound-state processes, the amplitude for a transition is not gauge-invariant. The rate is, however, because the positions of poles in Green's functions are gauge invariant and the rate is just the imaginary part (the energy level is the real part) of the pole.

A complete discussion, starting from first principles, of bound-state and purely atomic physics effects for the tree-level graph of Fig. 2 has already been carried out for atoms with an arbitrary number of electrons.[27] It is also possible to reformulate field theory in terms of QED bound-states, rather than the S-matrix, and to do a systematic perturbation expansion and renormalization around bound-states.[28] One can imagine, however, that the non-Abellian nature of the WS theory introduces all sorts of subtleties[29] and it remains to carry out a systematic calculation of all radiative corrections to PV effects in one-electron atoms within this context.

The experimental situation for PV effects in one-electron atoms seems quite hopeful. Groups at Yale, Univ. of Washington and Univ. of Michigan are working on experiments which will measure C^T_{2P} reliably at the tree level.[1] The progress in trapped ions reported by Prof. Toschek is so impressive that one can imagine experiments to measure PV effects in, say, trapped He^+. ^2H deuterium may provide a laboratory in which we might better understand strong-interaction effects because the nuclear-spin dependent $C^T_{2;^2H} = C^T_{2P} + C^T_{2N} = 0$ at tree level. It may even be possible to find some coherent PV effect in Bose-Einstein condensed spin-polarized hydrogen. As reported by Prof. Kleppner, spin-polarized hydrogen is already available at high densities and they are stalking the elusive Bose-Einstein gas relentlessly.

One of the most hopeful suggestions, to my mind, for the measurement of the radiative corrections to PV effects in atoms, is a comment by Dunford, Lewis and Williams.[5] They point out that any experiment which sums over hyperfine states in hydrogen and is carried out at the $\beta - f$ crossing in the hydrogen Rabi diagram at a constant external magnetic field $|\underset{\sim}{B}| \sim 1200$ Gauss will measure $C_{1P} + .105C_{2P}$. The reason for the suppression of C_{2P} is that, for $|\underset{\sim}{B}| \gtrsim 100$ Gauss, we are in the Paschen-Back limit and the atomic states have almost pure I_z (nuclear spin) and J_z (electron total angular momentum) rather than total angular momentum F_z; $\underset{\sim}{F} = \underset{\sim}{I} + \underset{\sim}{J}$. Ideally, such an experiment would be done with deuterium ^2H because it would measure, to extremely good accuracy, the radiative corrections to the combination $C_{1;^2H} = C_{1P} + C_{1N}$ and suppress the uncertainties in $C_{2P} + C_{2N}$ due to strong interactions. We have seen that C_{1P} and C_{1N} do not suffer terribly from strong-interaction uncertainties and so these experiments in hydrogen or deuterium will be theoretically 'clean' once the bound-state effects are understood. They will then test the radiative corrections in $SU(2)_L \times U(1)$. Since radiative corrections are incalculable in the phenomenological Fermi theory, this will provide the first evidence that non-Abellian gauge invariance, responsible for the deeper renormalizable SBGT structure of the Weinberg-Salam theory, is a fundamental principle of nature.

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(6) G. 't Hooft & M. Veltman, in "Particle Interactions at Very High Energies", edited by D. Speiser, F. Halzin and J. Weyers (Plenum, New York 1974), Part B, pg. 177; G. Passarino and M. Veltman, Nuc. Phys. B 160 (1979) 151; M. Consoli, Nuc. Phys. B 160 (1979) 208. These two papers develop quite an efficient machinery for the calculation of higher-order processes in $SU(2)_L X U(1)$.

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(8) The P $\overline{P}$ collider at CERN and, eventually, LEP there advertise that they will be able to measure these masses to better than 1%. Fermilab reports similarly hopeful experimental endeavors. If this is true, the masses M_W and M_Z will provide the cleanest (most free from strong interaction effects) data with which to calculate radiative corrections.

(9) This definition of s_θ^2 is originally due to A. Sirlin, Phys. Rev. D22 (1980) 971 and, we feel, gives the simplest renormalization scheme for $SU(2)_L X U(1)$. M. Gluck & E. Reya, Phys. Rev. Lett. 47 (1981) 1104 have shown that precise meaurements of s_θ^2 as defined by the ratio of charged to neutral current neutrino scattering (the usual definition) are hampered by strong-interaction effects at present energies. Notably, higher twist effects are shown to be larger than electro-weak corrections to certain neutrino processes and there results a 5% limit of accuracy in knowing s_θ^2. On the other hand, if we use M_W and M_Z as input parameters, the effects of strong interactions on these numbers can be reliably computed as they are evaluated at $|q^2| \simeq M_W^2, M_Z^2 >> 1 \text{ GeV}^2$, the scale of strong interactions and well after the onset of asymptotic freedom. This is one reason why we believe this renormalization scheme with $s_0^2 = 1 - M_W^2/M_Z^2$ will eventually be shown to be superior. For a general theorum on the large M_W, M_Z dependence of electroweak processes see A. Sirlin, Nuc. Phys. B 196 (1982) 83.

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Table I

Numerical Values of the Parity-Violation Parameters (30)

C_{1P}

s_0^2	M_Z(GeV)	Higgs' Mass M_H: 3.16 GeV	10 GeV	100 GeV	1 TeV	10 TeV	100 TeV
.175	91.5	.181	.181	.182	.189	.194	.200
.186	92.1	.145	.145	.147	.152	.157	.163
.197	92.7	.113	.113	.115	.119	.124	.130
.208	93.4	.0843	.0845	.0862	.0904	.0951	.100
.214	93.7	.0711	.0713	.0728	.0769	.0815	.0866
.219	94.1	.0584	.0586	.0601	.0641	.0686	.0736
.231	94.7	.0349	.0351	.0365	.0403	.0446	.0494
.242	95.4	.0135	.0137	.0150	.0186	.0227	.0273
.253	96.1	-.00610	-.00597	-.00477	-.00127	.00271	.00713
.264	96.9	-.0241	-.0240	-.0229	-.0195	-.0157	-.0114
.275	97.6	-.0408	-.0407	-.0397	-.0364	-.0327	-.0285

C_{1N}

s_0^2	M_Z(GeV)	Higgs' Mass M_H: 3.16 GeV	10 GeV	100 GeV	1 TeV	10 TeV	100 TeV
.175	91.5	-.556	-.556	-.559	-.564	-.569	-.575
.186	92.1	-.523	-.524	-.526	-.531	-.536	-.541
.197	92.7	-.494	-.495	-.497	-.501	-.506	-.511
.208	93.4	-.469	-.469	-.471	-.475	-.479	-.484
.214	93.7	-.457	-.457	-.459	-.463	-.467	-.472
.219	94.1	-.446	-.446	-.448	-.452	-.456	-.460
.231	94.7	-.425	-.425	-.427	-.431	-.434	-.438
.242	95.4	-.406	-.407	-.408	-.412	-.415	-.419
.253	96.1	-.389	-.390	-.391	-.395	-.398	-.401
.264	96.9	-.374	-.374	-.376	-.379	-.382	-.385
.275	97.6	-.360	-.360	-.362	-.364	-.367	-.370

Table I (continued)

C_{2P}

s_0^2	M_Z(GeV)	Higgs' Mass M_H: 3.16 GeV	10 GeV	100 GeV	1 TeV	10 TeV	100 TeV
.175	91.5	.214	.214	.216	.223	.230	.238
.186	92.1	.170	.170	.173	.179	.185	.193
.197	92.7	.131	.132	.134	.139	.146	.153
.208	93.4	.0968	.0970	.0990	.104	.110	.117
.214	93.7	.0808	.0811	.0830	.0882	.0939	.100
.219	94.1	.0657	.0659	.0678	.0728	.0784	.0847
.231	94.7	.0376	.0378	.0395	.0443	.0497	.0557
.242	95.4	.0120	.0122	.0138	.0184	.0236	.0294
.253	96.1	-.0113	-.0111	-.00959	-.00520	-.000204	.00534
.264	96.9	-.0326	-.0325	-.0311	-.0268	-.0220	-.0166
.275	97.6	-.0523	-.0522	-.0508	-.0468	-.0421	-.0368

C_{2N}

s_0^2	M (GeV)	Higgs' Mass M_H: 3.16 GeV	10 GeV	100 GeV	1 TeV	10 TeV	100 TeV
.175	91.5	-.208	-.209	-.211	-.218	-.225	-.232
.186	92.1	-.165	-.165	-.168	-.174	-.180	-.188
.197	92.7	-.127	-.127	-.129	-.135	-.141	-.148
.208	93.4	-.0926	-.0928	-.0949	-.100	-.106	-.113
.214	93.7	-.0768	-.0770	-.0790	-.0841	-.0899	-.0963
.219	94.1	-.0618	-.0620	-.0639	-.0689	-.0746	-.0808
.231	94.7	-.0340	-.0342	-.0359	-.0407	-.0461	-.0521
.242	95.4	-.00873	-.00891	-.0105	-.0151	-.0203	-.0260
.253	96.1	.0143	.0141	.0126	.00825	.00325	-.00229
.264	96.9	.0354	.0353	.0339	.0296	.0248	.0194
.275	97.6	.0548	.0547	.0534	.0493	.0446	.0394

Table II

Ratio of One-Loop Corrections to Tree Values of Parity-Violation Parameters for $s_\theta^2 = .23$

M_H	C^1_{1P}/C^T_{1P}	C^1_{1N}/C^T_{1N}	C^1_{2P}/C^T_{2P}	C^1_{2N}/C^T_{2N}
3.16 GeV	.0294	-.0265	-.118	-.202
10 GeV	.0341	-.0257	-.113	-.197
100 GeV	.0750	-.0211	-.0723	-.157
1 TeV	.187	-.0130	.0400	-.0443
10 TeV	.314	-.00487	.167	.0823
100 TeV	.454	.00432	.307	.223

Figure Captions

Figure 1a: Lowest Order PV electron-quark scattering in Fermi Theory
Figure 1b: $\mathcal{O}(G_F^2)$ Correction to electron-quark scattering in Fermi Theory

Figure 2 : Lowest-Order PV electron-quark scattering in Weinberg-Salam Theory

Figure 3 : Feynman Rules Generated from $L^b_{\psi W}$

Figure 4 : Possible topologies of Feynman Rules generated from L^b

Figure 5a: 1PI electron -photon Vertex Part
Figure 5b: 1PI Z-A mixing Part
Figure 5c: 1PI Electron Self-Energy Part

Figure 6 : Feynman Diagrams 1P-reducible by cutting a Z line

Figure 7 : Feynman Diagrams 1P-reducible by cutting a photon line

Figure 8 : Box Diagrams

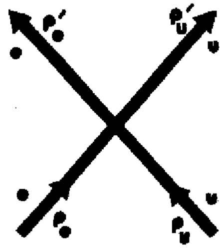


FIG 1a

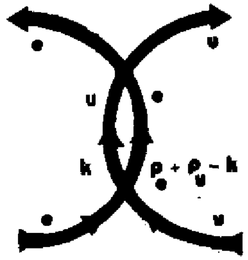
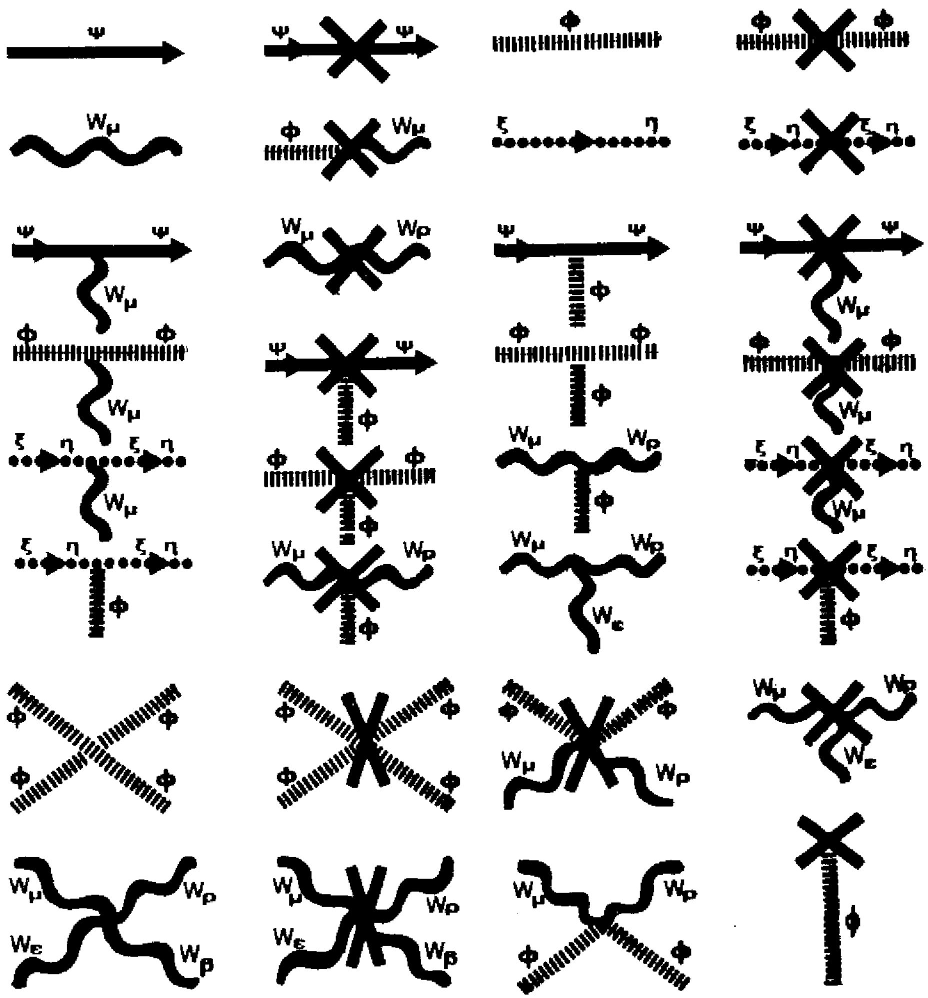


FIG 1b

FIG 4



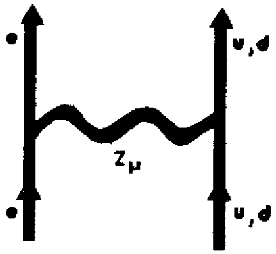
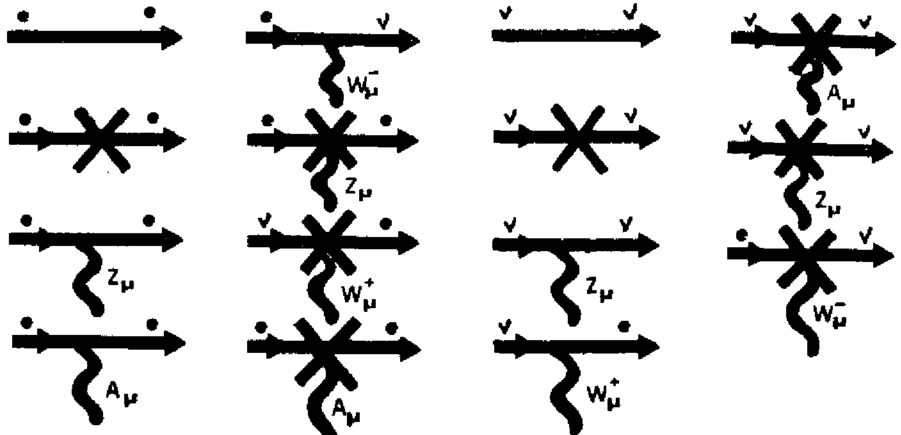


FIG. 2



FIG

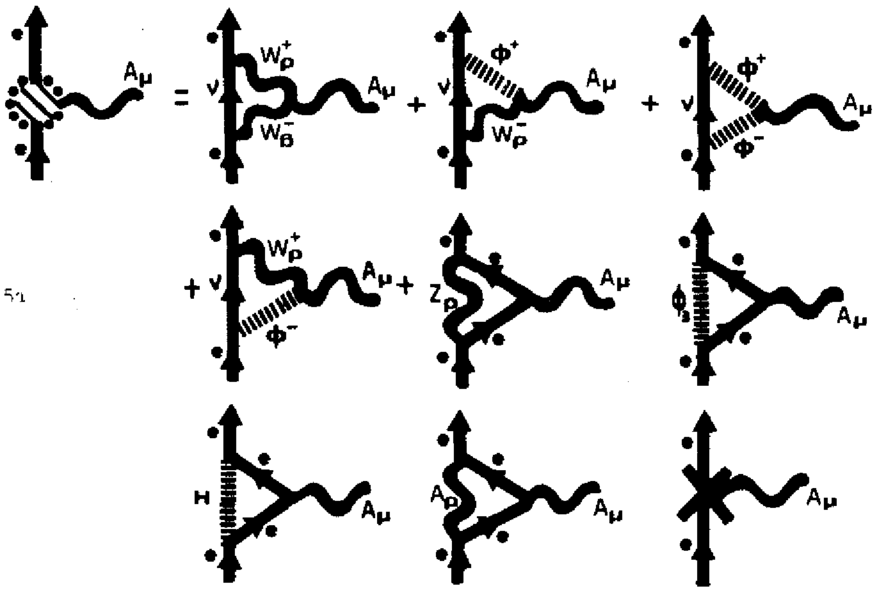


FIG 5a

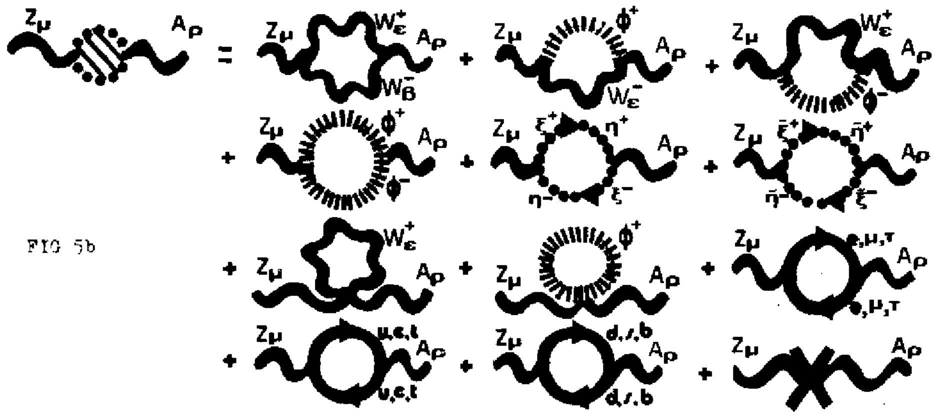


FIG 5b

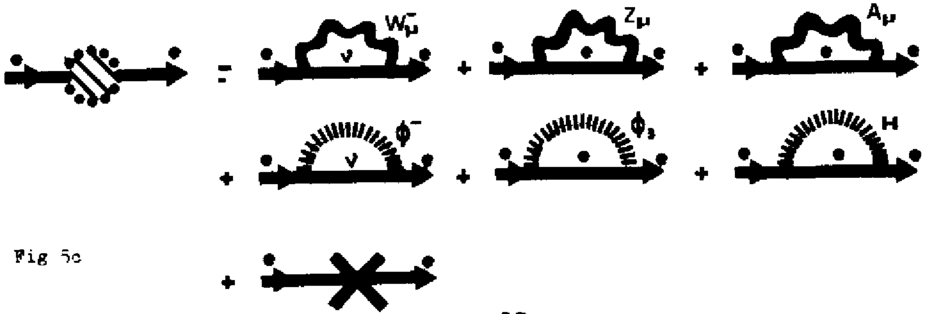


Fig 5c

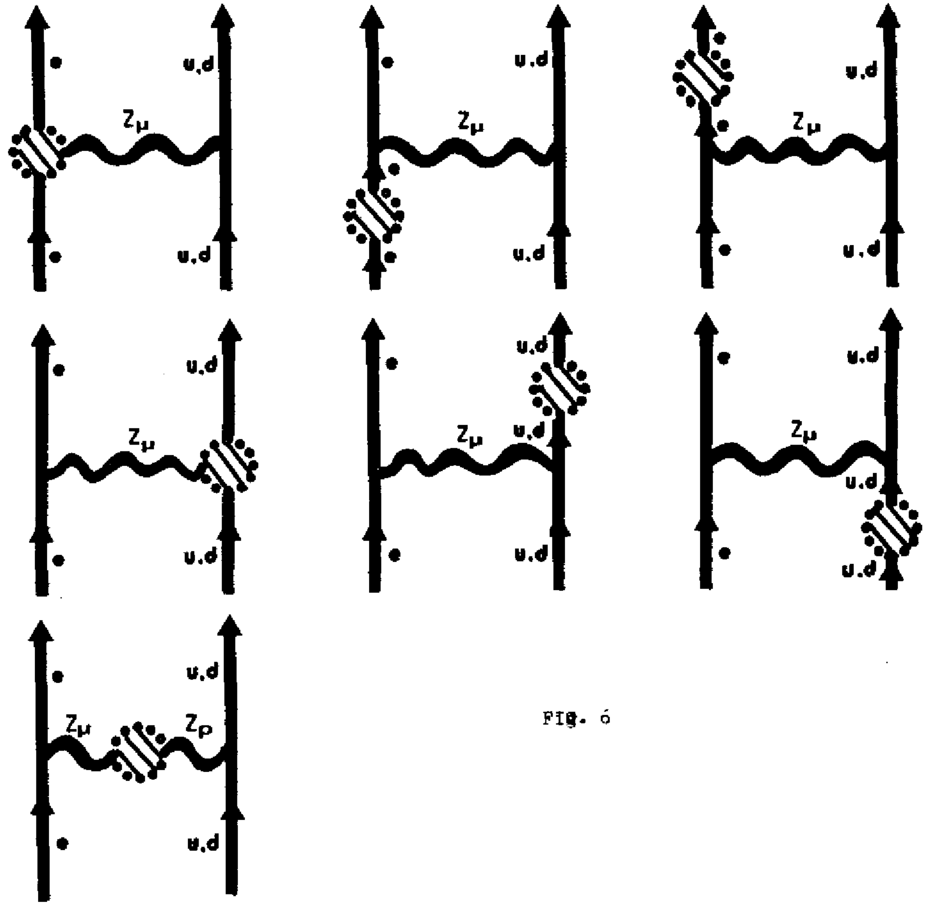


FIG. 6

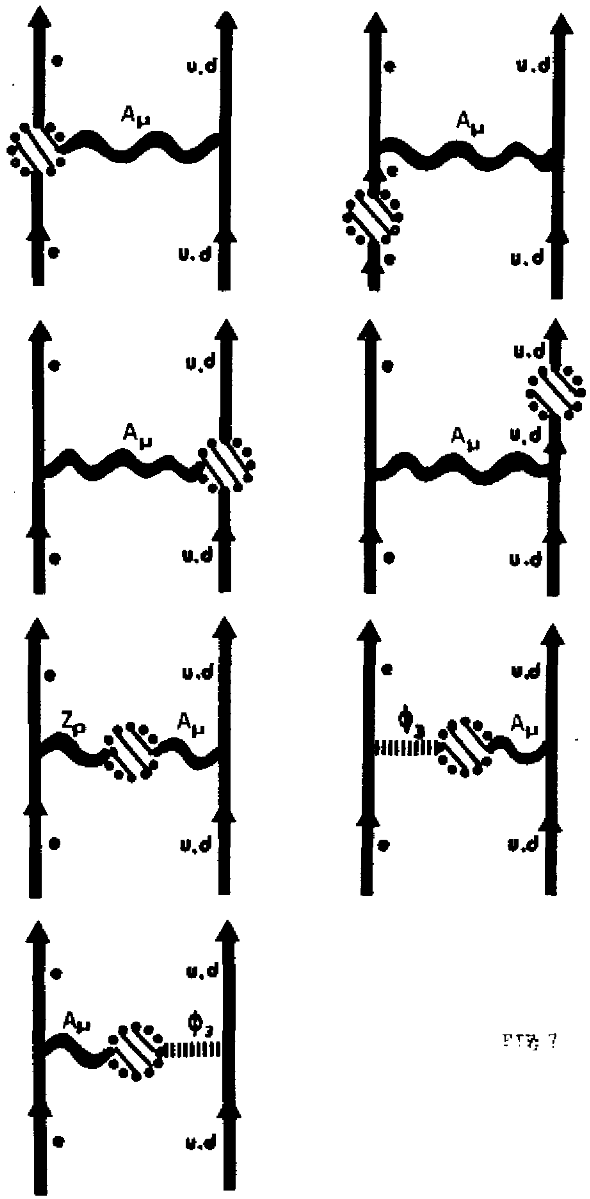


FIG 7

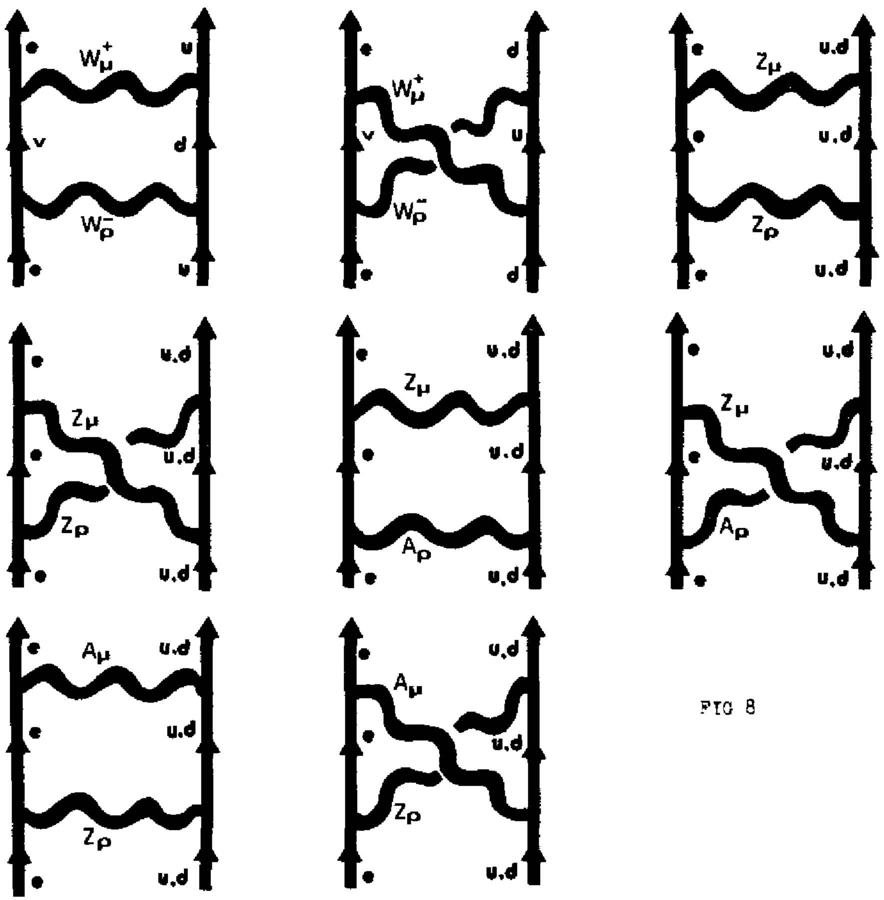


FIG 8

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