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# O(4) Treatment of the Electromagnetic-Weak Synthesis*

A. Pais
*Rockefeller University, New York, New York 10021*
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A gauge theory is outlined in which the existence of the Cabibbo angle $\theta$, leptonic $CP$-invariance violation, and $\mu$-$e$ nonuniversality are linked. To $O(G)$, the sole source of nonleptonic decays is a neutral current. There is no conflict with present neutrino data. This $CP$-invariance violation induces only superweak effects in $K$ decays. A strategy for finding $\theta$ is indicated.

Imagine a semileptonic interaction of the form[1] $ig_1(\overline{\mathcal{P}}\mathcal{N}+\overline{\nu}_e e+\overline{\nu}_\mu \nu)_L W^1+ig_2\{\overline{\mathcal{P}}\lambda+a(e^{i\varphi}\overline{\nu}_e e+e^{i\psi}\overline{\nu}_\mu \mu)\}_L W^2$ + H.c. $W^{1,2}$ are charged vector bosons with masses $M_1\neq M_2$; $\mathcal{P}$, $\mathcal{N}$, and $\lambda$ are quarks. The real number $a$ and the phases $\varphi$ and $\psi$ are observable. To leading order, $\mu e$ universality is valid in $\mathcal{N}\to\mathcal{P}$ and in $\lambda\to\mathcal{P}$ $\beta$ decays. The condition that the absolute values of the amplitudes for $\mu\to e$, $\mathcal{N}\to\mathcal{P}$, $\lambda\to\mathcal{P}$ are in the ratios[1] $1:\cos\theta:\sin\theta$ implies $(1-a^2)\tan\theta=2a\cos(\psi-\varphi)$. If $\psi=\varphi$, $a\neq 1$, $\mu e$ universality is strict. If $\psi\neq\varphi$, it is violated together with $CP$. To $O(g_1^2; g_2^2)$, $CP$ and $T$ are conserved as is seen by redefining phases; but not, in general, in higher order. In this note, these ideas are examined[2] in the context of a $CPT$-invariant O(4)-gauge theory, with the choice $a=1$, $\varphi=0$, $\psi=\pi/2$. While then $CP$-invariance violation is maximal for the muonic terms in the $\Delta Q=\pm 1$ currents, the physical effects thereof will turn out to be minuscule for[1] $q^2/M^2\ll 1$.

Six gauge fields, $A_\mu{}^i$, $C_\mu{}^i$, $i=1, 2, 3$, appear in the gauge-invariant derivative $D_\mu=\partial_\mu-i(g_1\vec{A}_\mu\cdot\vec{t}+g_2\vec{C}_\mu\cdot\vec{\rho})$. Here $[\vec{t},\vec{\rho}]=0$; $\vec{t}\times\vec{t}=i\vec{t}$; $\vec{\rho}\times\vec{\rho}=i\vec{\rho}$. The charge operator is $eQ$; $Q=t_3+\rho_3$. For now, we bypass the option $g_1/g_2\neq 1$ and put

$$g_1=g_2=e\sqrt{2}. \quad (1)$$

Then there is a further invariance under reflections $R$:$\vec{t}\leftrightarrow\vec{\rho}$ in O(4). Consider a scalar field quartet $H$ with charges (+,0,0,−). Here the action of $\vec{t},\vec{\rho}$ is $2\vec{t}=\vec{\tau}\otimes 1$; $2\vec{\rho}=1\otimes\vec{\tau}$. 1, $\vec{\tau}$ are Pauli matrices. Let $H$ have vacuum expectation values $\langle 0,a,a',0\rangle$, *a,a' real*. Denote such an $H$ as $H(a,a')$. Introduce two such $H$'s: $H(a,a')$ and $H(0,b)$; $a,a'$, $b\neq 0$. The $H$'s generate vector masses and $D_\mu$ becomes

$$D_\mu=\partial_\mu-ieQA_\mu-ie(t_3-\rho_3)Z_\mu-ie\cdot 2^{-1/2}\{W_\mu{}^1(t_+-\rho_+)+\text{H.c.}\}, \quad (2)$$

$$A_\mu\sqrt{2}=A_\mu{}^3+C_\mu{}^3, \quad Z_\mu\sqrt{2}=A_\mu{}^3-C_\mu{}^3, \quad (3)$$

$$2W_\mu{}^{1,2}=[A^1\mp C^1-i(A^2\mp C^2)]_\mu, \quad \xi\equiv M_0{}^2/(M_1{}^2+M_2{}^2)=1, \quad (4)$$

where $M_0$ is the mass of the neutral heavy $Z$-vector meson. $W_\pm^2$ and $A$ are even under $R$ and are associated with a subalgebra O(3); $W_\pm^1$ and $Z$ are $R$ odd.

All[1] $f_L$ are grouped in quartets $(\frac{1}{2},\frac{1}{2})$. The electron (muon) quartet will be denoted by $E_L$ ($M_L$). For baryons, two quartets $Q_L{}^1$, $Q_L{}^2$ of quarks will be introduced. If we assume *no group extension* (see the final comment below) then we need $E_R$, $M_R$, $Q_R{}^1$, $Q_R{}^2$, all triplets. The option (1,0) or (0,1) affects the fermion mass problem, but to the orders concerned makes no difference in the phenomenology.[3] Let us say for the moment that the triplets are (1,0). In addition $n_R$, $\lambda_R$ will be singlets. The definitions are [with $\epsilon = \exp(i\pi/4)$]

$$E_L = (x^+, x^0, -\nu_e, -e)_L, \quad M_L = (y^+, \epsilon\cdot 2^{-1/2}[\nu + y^0], -\epsilon^*\cdot 2^{-1/2}[\nu - y^0], -\mu)_L,$$
$$Q_L{}^1 = (\mathcal{P}, \tfrac{1}{2}[\mathfrak{N} + \lambda + q^0\sqrt{2}\,], \tfrac{1}{2}[\mathfrak{N} - \lambda + r^0\sqrt{2}\,], -q^-)_L, \quad Q_L{}^1 = (q^+, \tfrac{1}{2}[\mathfrak{N} - \lambda - r^0\sqrt{2}\,], -\tfrac{1}{2}[\mathfrak{N} + \lambda - q^0\sqrt{2}\,], -r^-)_L,$$
$$E_R = \text{“}(x^+, x^0, e)\text{”}_R;\; M_R = \text{“}(y^+, y^0, \mu)\text{”}_R;\; Q_R{}^1 = \text{“}(\mathcal{P}, r^0, r^-)\text{”}_R;\; Q_R{}^2 = \text{“}(q^+, q^0, q^-)\text{”}_R. \tag{5}$$

The quotation marks on the $R$ triplets indicate that the particle symbols are defined up to a phase (about which more later). Thus we postulate two heavy leptons of $E$ type, two of $M$ type. As to the $Q$'s, $\mathcal{P}$, $\mathfrak{N}$, $\lambda$ are conventional. The other five particles are supposed to carry (one or more) further quantum numbers, collectively denoted as $C$. At least one of these is additive. It is left open at this time whether they carry a nonzero $I$, $Y$ as well.

$J^{(0,1,2)}$ will denote the currents coupled to $Z$, $W^1$, $W^2$. $\mu$ decay via $J^{(1,2)}$ sets the scale to be (if there is no group extension)

$$8G = e^2(2M_1{}^{-4} + 2M_2{}^{-4})^{1/2};\quad M_1 = 37.2\cos^{-1/2}\theta,\quad M_2 = 37.2\sin^{-1/2}\theta, \tag{6}$$

but let us consider $J^{(0)}$ first:

$$J^{(0)} = -ie[\{\bar{\nu}_e\nu_e - \bar{x}^0x^0 - \bar{\nu}_\mu y^0 - \bar{y}^0\nu_\mu - ((\bar{\mathfrak{N}} + \bar{\lambda})q^0 - (\bar{\mathfrak{N}} - \bar{\lambda})r^0 + \text{H.c.})/\sqrt{2}\}_L + \textstyle\sum (f^+f^+ - f^-f^-)_R], \tag{7}$$

where $\sum$ denotes the sum over all fermions occuring in all the $R$ triplets of Eq. (5). Thus, within the present approach, $J^{(0)}$ exhibits a "maximal" $\mu e$ nonuniversality. The option to put the $CP$-invariance violation in $M$ rather than in $E$ has been chosen to avoid conflict with available high-energy $\nu$ data.[4] Up to small $\nu_e$ impurities,[5] these experiments should indeed yield null results if there exists a $J^{(0)}$ of the structure Eq. (7). For elastic $\bar{\nu}_e e$ scattering one finds, with $x = \sin 2\theta$,

$$\sigma(\bar{\nu}_e e) = \frac{4}{3}\frac{G^2 m E_\nu}{\pi}\frac{1 + 2x + x^2(1 + 3\xi^{-2})}{1 + x} \simeq 0.95\times 10^{-41}E_\nu \text{ cm}^2/\text{electron}, \tag{8}$$

with $\xi = 1$ [Eq. (4)] and $x \simeq 0.4$. This is $\simeq 1.7$ times the $V-A$ value, well within experimental limits.[6] From the present point of view $\sigma(\bar{\nu}_e e)$ is of interest as an independent measurement of the Cabibbo angle. To obtain $\sigma(\bar{\nu}_e e)$, multiply the right-hand side of Eq. (8) by 3 and put $3\xi^{-2} \to \xi^{-2}/3$. Up to factors $\sim 1$, $\nu_\mu e$ scattering is as in Fujikawa *et al.*[7]

Turning to the $Q$'s in Eq. (7), to $O(e^2)$ an effective nonleptonic interaction $\sim G(\bar{\lambda}\mathfrak{N})_L(\bar{q}^0q^0 - \bar{r}^0r^0)_L + \text{H.c.}$ is generated. It satisfies $\Delta C = 0$, $|\Delta S| = 1$. We may now follow a line of reasoning due to Lee and Treiman,[8] who were the first to recognize precisely this kind of four-fermion interaction in a different context: Considering the usual mesons and baryons as the conventional states but with $\Delta C = 0$ pairs, this interaction projects onto those usual states and generates nonleptonic decays. These are $\Delta I = \frac{1}{2}$ if $(\bar{q}^0q^0 - \bar{r}^0r^0)_L$ is a $\Delta I = 0$ operator (as is true, for example, if $q^0$, $r^0$ are $I = 0$). A more detailed discussion must await the embedding of this scheme in hadronic symmetries.

Consider next the contributions of the $L$ leptons:

$$\begin{aligned}&\text{to } J^{(1)},\quad -ie[\bar{x}^+(x^0 + \nu_e) + (\bar{\nu}_e + \bar{x}^0)e + \bar{y}^+(\nu_\mu + iy^0) + (\bar{\nu}_\mu - i\bar{y}^0)\mu]_L/\sqrt{2}\,;\\ &\text{to } J^{(2)},\quad -ie[\bar{x}^+(x^0 - \nu_e) + (\bar{\nu}_e - \bar{x}^0)e + \bar{y}^+(y^0 + i\nu_\mu) + (i\bar{\nu}_\mu - \bar{y}^0)\mu]_L/\sqrt{2}\,.\end{aligned} \tag{9}$$

Heavy-lepton decay violates universality. To $O(e^4)$, $W^1$ and $W^2$ mix via $E$ bubbles and $M$ bubbles. For $q^2$ values of the $W$'s $\ll M^2$ the off-diagonal mass is given by

$$W_{12} = \frac{-\alpha}{2\pi}\left[\frac{m^2(x^+)m^2(x^0)}{m^2(x^+) - m^2(x^0)}\ln\frac{m^2(x^+)}{m^2(x^0)} + i\,m^2(y^0)\ln\Lambda^2\right]\frac{1}{M^2}, \tag{10}$$

and $W_{21} = W_{12}{}^*$. $\Lambda$ is a cutoff (in heavy $M$-lepton mass units). The one renormalization needed will be discussed more fully elsewhere. We make the common assumption that this does not change the order of magnitude of the effect. To $O(G\alpha)$, $W_{12}$ or $W_{21}$ insertions in the propagator give only an irrelevant overall phase to the $\mu$-decay amplitude $a_\mu$. This is no longer true if we also include electromagnetic corrections. The $CP$-invariance violation in $a_\mu$ is therefore $O(G\alpha^2)$.

The other potential $CP$-invariance violator $O(iG\alpha)$ stems from the "box graph" due to $W_1Z$ exchange between an $M$ and an $E$ line. Such graphs have recently been the subject of detailed study.[7,9] Following with care the reasoning of Ref. 7, one sees again that only an overall phase results as long as $(q^2)_{\rm leptons} \ll M^2$. Again electromagnetic corrections can bring out the effect. Thus in low-momentum transfer lepton processes the amplitude effect is $\lesssim O(G\alpha^2)$. Note the usefulness of $CPR$ invariance for $M$ couplings: Where $CP$ is nonconserved, so is $R$, but $CPR$ is conserved, as in the $M$ contribution to $W_{12}$ and $W_1Z$ exchange. As a further example, all triangle insertions in the (anomaly-free) $W_1W_2Z$ vertex conserve $CP$. The $R$ leptons [since they are (1,0) couple equally to $W^1$ and $W^2$. Their presence is easily shown not to affect the $CP$ estimates just made.

The $Q_L$'s contribute

$$\begin{aligned} &\text{to } J^{(1)}, \quad -ie[\bar{\mathcal{P}}(\mathfrak{N}+A_1^0)+(\bar{\mathfrak{N}}+A_1^0)q^-+\bar{q}^+(\mathfrak{N}-A_1^0)+(\bar{\mathfrak{N}}-\bar{A}_1^0)r^-]_L\sqrt{2}\,; \\ &\text{to } J^{(2)}, \quad -ie[\bar{\mathcal{P}}(\lambda+A_2^0)-(\bar{\lambda}+\bar{A}_2^0)q^- -\bar{q}^+(\lambda-A_2^0)+(\bar{\lambda}-\bar{A}_2^0)r^-]_L/\sqrt{2}\,, \end{aligned} \tag{11}$$

with the abbreviations $A_1^0=(q^0+r^0)/\sqrt{2}$ and $A_2 = (q^0-r^0)/\sqrt{2}$. Following the arguments of Ref. 9, Eqs. (7) and (11) show the following: (1) The amplitude for $\lambda\lambda\to\mathfrak{N}\mathfrak{N}$ is at most $O(G\alpha\Delta^2)$, independently of the relative values of $M_{0,1,2}$. Here $\Delta$ denotes generically[11] the ratio of ($Q$-mass)$^2$ differences to any of the $(M_{0,1,2})$. Making the common assumption[9] that $|\Delta|\ll 1$, $\Delta S=2$ is amply suppressed. (2) The same is true for $K_L\to\mu\bar{\mu}$ which is $O(iG\alpha\Delta m^2(y^0)M^{-2})$. The $Q_R$'s couple equally to $W^1$ and $W^2$; $\mathfrak{N}_R$ and $\lambda_R$ are (0,0) and do not appear in the currents, hence do not affect these estimates.

Can virtual lepton effects induce $CP$-invariance violation in $K$ decays of the desired order of magnitude, via $W_{12}$ effects for example? Following nearly verbatim the $\mu$-decay argument one readily concludes that the on-shell effects in $\lambda\to\mathcal{P}e\bar{\nu}_e$ or $\lambda\to\mathcal{P}\mu\bar{\nu}_\mu$ (hence in $K_{l2}, K_{l3}$) are at most $O(G\alpha^2)$, one factor $\alpha$ too small. Lepton loops also induce "$Q$ mixings" such as $\lambda\leftrightarrow\mathfrak{N}$, for example. Writing Eq. (10) as $W_{12}=(\alpha/\pi)(\omega_e+i\omega_\mu)$, such insertions are $O((\alpha/\pi)\omega_e\Delta m_Q)$ for $\lambda\leftrightarrow\mathfrak{N}$ via an $E$ loop; via an $M$ loop, put $\omega_e\to\pm i\omega_\mu$ for $\lambda\leftrightarrow\mathfrak{N}$. Here $m_Q$ is a typical $Q$ mass and the crude estimate holds for virtual $Q$'s with $p^2\ll m_Q^2$. Such insertions also turn out to be too small, given the constraints imposed by (real) $\lambda\lambda\to\mathfrak{N}\mathfrak{N}$ and $K_L\to\mu\bar{\mu}$. Hence the impact of the leptonic $CP$-invariance violation on $K_L$ decays is superweak.

I have no proof that the present mechanism can yield the requisite superweak magnitude $\delta m_{im} \approx 6\times10^{-15}i$ MeV for the imaginary part of the off-diagonal mass matrix element $K_0\to\bar{K}_0$ but would like to give an extremely crude argument to indicate that such a magnitude at least does not appear to do violence to any of the typical parameters encountered: Consider the sequence $K^0 \to\bar{\lambda}\mathfrak{N}\to\lambda\bar{\mathfrak{N}}\to\bar{K}_0$ via a $Q\bar{Q}$ loop with one $\omega_e$, one $\omega_\mu$ insertion, and use a vertex of unit strength to symbolize a strong $(K^0\bar{\lambda}\mathfrak{N})$ coupling. Then $\delta m_{im} \sim i(\alpha/\pi)^4\Delta^2\omega_e\omega_\mu m_Q^2 m_K^{-1}$. To give just one example, this gives reasonable results if $m_Q\sim 10$ GeV, and $\Delta, \omega_e, \omega_\mu$ all crudely $\sim\alpha$. Such orders do not seem unfair and reflect the typical feature of gauge theories that factors $\alpha$ can appear via mass ratios.

There remain the crucial and profound fermion mass problems. Minimally, one has to show that all matrices can be diagonalized, and in such a way that no degeneracies remain which would render the foregoing results invalid. In addition, the following questions must be faced. (1) There is an embedding problem: to find a match between spontaneously broken local symmetries and the broken hadron symmetries[11] through mass formulas, etc. (2) If $\mu e$ nonuniversality is built into a scheme, then [regardless of whether the particular O(4) realization attempted here is correct or not] one will ask for clues to distinct mass-generating mechanisms for $\mu$ and $e$. (3) From constraints on baryon masses or on lepton masses (or both) the value of $\theta$ should eventually be fixed. A few comments on this program will be presented here.

(1) *The E masses.*—Consider the couplings $\bar{E}_L\{\vec{u}\cdot\vec{E}_R H(a,a')+\vec{v}\cdot\vec{E}_R H(0,b)\}$ + H.c. which set the scale of $a,a',b$. We have the options[12] $(\vec{u},\vec{v})=(\vec{t},\vec{t})$, $(\vec{\rho},\vec{\rho})$, $(\vec{t},\vec{\rho})$, or $(\vec{\rho},\vec{t})$. $(\vec{t},\vec{t})$ diagonalizes provided $a'=-b$, and yields

$$m(e)=0, \quad m(x^+)=m(x^0)\sqrt{2}\,. \tag{12}$$

$(\vec{\rho},\vec{\rho})$ interchanges $x^+\leftrightarrow e$ in Eq. (12). $(\vec{\rho},\vec{t})$ is remarkable. It yields $a'=-b$, $m(x^+)+m(e)=m(x^0)\sqrt{2}$. Let us *impose* $m(e)=0$. Then Eq. (12) reappears but the extra constraint implies $a=a'=-b$ so that[3]

$$\tan\theta=\frac{b^2+(a-a')^2}{b^2+(a+a')^2}\to\tfrac{1}{5}, \tag{13}$$

which yields the attractive "bare" value $\theta = 11°20'$. This is far from saying that $\theta$ has been computed, but rather to provide an example of how $\theta$ may eventually be fixed by mass constraints. Either coupling scheme just mentioned fixes the phases in $E_R$ [see Eq. (5)] to $(x^+, -x^0, -e)$. Such alternatives will be discussed further elsewhere.

(2) *The M masses.*—Consider $M_L\{A\vec{u}\cdot\vec{M}_R H(a,a') + B\vec{v}\cdot\vec{M}_R H(0,b)\}$ + H.c. $A, B$ may be relatively complex! As an example, retain $a = a' = -b$ just encountered and take $(\vec{u}, \vec{v}) = (\vec{t}, \vec{p})$. Then diagonalization is achieved for $M_R = (-iy^+, -\epsilon y^0, +\mu)_R$; $A\sqrt{2} = B\epsilon$ with the following zeroth-order mass relation: $m(y^+) = m(y^0) = m(\mu)$. Thus a constraint imposed to put the bare $m_e = 0$ can yield a nonzero bare average mass for the $M$ multiplet. One is now at a crossroads which is being examined further. Either big self-energy effects split the multiplet, or new Higgs fields must be introduced. The latter is possible without annihilating Eq. (3). Indeed any further Higgs field which is a representation of O(4) and reduced with respect to $R$ maintains Eq. (3). I have shown that one of the remaining options for such fields, $(\frac{3}{2}, \frac{1}{2}) + (\frac{1}{2}, \frac{3}{2})$ and $(\frac{3}{2}, \frac{1}{2}) - (\frac{1}{2}, \frac{3}{2})$, suffices to split the masses further. Two quantities are affected: the $\xi$ parameter, Eq. (4), and the expression, Eq. (13), for $\tan\theta$. The question of the possible approximate validity of Eqs. (4) and (13) poses new constraints, which demand a further examination of the Higgs potential surface.[13] It is not the purpose to pursue these questions at this place, but merely to furnish a few examples of sequential constraints which arise in problems of this kind.

(3) *Q masses.*—It can be shown that the totality of Higgs fields described above can give nondegenerate masses to all particles. Had one chosen $Q_R$'s to be $(\frac{1}{2}, \frac{1}{2})$ then it can be shown that massless particles remain in the zeroth order approximation.

As a final comment, the secondary role of the $f_R$ in the phenomenology raises the question whether one is forced to assign them as done here, or whether one can proceed otherwise. A study of O(4)⊗ U(1) and of O(4)⊗O(4) shows that all the promising properties of $J^{(0,1,2)}$ can be maintained in a description where some or all of the $f_R$ are scalar with respect to O(4). Up to scale[14] the structure of $J^{(0,1,2)}$ remains unaffected, except that some or all of the $\bar{f}_R f_R$ terms may no longer appear in $J^{(0,1,2)}$. It may perhaps be of use to keep such options in mind for the mass problems and for the understanding of connections with strong interaction dynamics.

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[1]Notations: $(\bar{\mathcal{P}}\mathfrak{N})_L \equiv \bar{\mathcal{P}}_L\gamma_\mu\mathfrak{N}_L$; $\mathfrak{N}_L \equiv (1+\gamma_5)\mathfrak{N}/2$; similarly for $R$ terms. $W^{1,2}$ is short for $W_\mu^{\ 1,2}$. $\theta$ is the absolute value of the Cabibbo angle. Mass units are GeV/$c^2$; $M$ will denote a mass $O(M_1) = O(M_2)$. $q^2$ = (momentum transfer)$^2$; $f_L, f_R$ stands for any left- (right-) handed fermion. Conventional electromagnetic and strong $C, P, T$ properties are assumed throughout.

[2]For an excellent review of gauge theories see B. W. Lee, in Proceedings of the Sixteenth International Conference of High Energy Physics, National Accelerator Laboratory, Batavia, Illinois, 1972 (to be published).

[3]In the phenomenology, only the signs of the $(\bar{f}f)_R$ terms in $J^{(0)}$, Eq. (7) are at stake. Of all effects considered, the only thing affected is a term (neglected anyway) $\sim m_e$ in $\sigma(\bar{\nu}_e e \to \bar{\nu}_e e)$. This is one of several options related to $R$ invariance. Thus, taking $W^1$ ($W^2$) to be $R$-even (-odd) changes $\theta$ to $\pi/2 - \theta$.

[4]W. Y. Lee, Phys. Lett. 40B, 423 (1972); D. H. Perkins, Oxford University Report No. 65/72, 1972 (to be published).

[5]$\sim 3\%$ in the Columbia experiment, R. Burns *et al.*, Phys. Rev. Lett. 15, 541 (1965); $\lesssim 0.6\%$ in the CERN experiment, G. Kalbfleisch, Nucl. Phys. B25, 197 (1970).

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[7]K. Fujikawa, B. W. Lee, A. I. Sanda, and S. B. Treiman, Phys. Rev. Lett. 29, 682, 823(E) (1972).

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[9]B. W. Lee, J. R. Primack, and S. B. Treiman, Phys. Rev. D (to be published).

[10]Three independent ratios appear: $m^2(p) - m^2(q^+)$; $m^2(q^-) - m^2(r^-)$, $m^2(q^0) - m^2(r^0)$, all divided by an $M^2$.

[11]S. Weinberg, Phys. Rev. Lett. 29, 388 (1972). Cf. also H. Georgi and S. Glashow, Phys. Rev. D 6, 2977 (1972).

[12]At this point, $R$ invariance can serve to select options, in regard to the behavior under reflections of the fermion triplets as well as of the Higgs quartets. $E_R = (1, 0)$ and $E_R = (0, 1)$ need separate consideration as will be shown elsewhere.

[13]This was especially emphasized to me by B. W. Lee.

[14]As compared to Eqs. (1), and (6), $g_1$, $g_2$, $M_1$, $M_2$ increase by a factor $(\frac{3}{2})^{1/2}$ for O(4)⊗ U(1), and by 2 for O(4)⊗ O(4). $\xi < 1$ in either case. In some such schemes it is possible for an approximate $\gamma_5$ symmetry to give zero for the $O(\alpha)$ electron self-energy if bare $m(e) = 0$; cf. S. Weinberg, Ref. 11.