Weak-Interaction Model with Finite Self-Masses of Leptons

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We propose a model of the weak interactions for leptons, including scalar intermediate bosons $C^{\pm}$ with a negative metric in addition to the usual weak vector bosons $B^{\pm}$. With this modification, problems of divergence and high-energy behavior are greatly reduced. If the logarithmic weak and electromagnetic self-mass divergences are assumed to cancel each other, the coupling constant g and the mass m_B $(=m_C)$ of the weak boson are predicted to be $g^2=\frac{3}{2}e^2$ and $m_B=137.7m_p$ $(\simeq m_p/\alpha)$ (where m_p is the proton mass), respectively. The finite self-masses are $\delta m_l\simeq(3e^2/16\pi^2)m_l \ln(m_B/m_l)^2$ $(\simeq 0.044m_e$ for an electron). It is shown that the renormalization of both the electromagnetic and weak interactions can be consistently accomplished in spite of the existence of the parity-violating interaction. The contributions of the weak interaction to the anomalous magnetic moments of leptons are calculated to be $\frac{1}{2}(g_l-2)=-Gm_l^2/12\sqrt{2}\pi^2$ $(\simeq -7.7\times10^{-10}$ for $\mu)$. Various weak reactions of leptons are discussed. For example, $\sigma(\nu_\mu+e\to\mu+\nu_e)$ approaches a constant value $G^2m_B^2/\pi\simeq2.8\times10^{-34}$ cm^2 at high energies.

I. INTRODUCTION

It has been known for a long time that the usual weak interaction, in which four fermion fields are locally coupled, has much worse properties from the field-theoretical point of view than the electromagnetic interaction. It behaves worse at high energies so that it may not be renormalizable. On the other hand, the divergence problem of lepton self-masses has been too difficult to be solved. In attempting to solve this problem, we think it useful to remember that leptons have two different types of interactions, electromagnetic and weak, and have no other interaction. We must, above all, do our best to find a possible solution of the problem, making use of only the above-mentioned fact.

In a previous Letter,[1] we proposed a model of the finite self-masses of leptons. The point of the model is as follows: We first consider a model of the weak interaction with scalar intermediate bosons $C^{\pm}$ *with a negative metric* in addition to the usual weak vector bosons $B^{\pm}$. With this modification, the effective Fermi interaction becomes better behaved at high energies, since the interaction Lagrangian has the form

$$\mathcal{L}_{\rm eff}=g^2\int d^4x' J_\lambda{}^+(x)\Delta_{\lambda\mu}(x-x')J_\mu(x'), \qquad (1.1)$$

where

$$\Delta_{\lambda\mu}(x-x')=[1/(2\pi)^4]\int d^4k \exp[ik(x-x')] \times[\delta_{\lambda\mu}/(k^2+m_B{}^2)]. \qquad (1.2)$$

Thus, the weak self-masses of the leptons come to have only logarithmic divergences. Second, we assume that these divergences should cancel those of electromagnetic self-masses of leptons. The purpose of this paper is to give details and other effects of the model.

Three different formalisms to give the same effective interaction (1.1) are presented in the Sec. II. In Sec. III, we obtain a prediction of the coupling constant g and the mass of the weak boson m_B from the second assumption. The method of renormalization in the model is shown in Sec. IV. In Secs. V and VI, we consider the effects on anomalous magnetic moments and various weak reactions of leptons, respectively. We use Sec. VII for discussion of the results, higher-order corrections, and generalizations of the model.

II. SCALAR BOSONS WITH NEGATIVE METRIC

We present three different formalisms to give the same effective weak interaction (1.1) but somewhat different electromagnetic interactions. All of them have an indefinite metric, η (Hermitian), of the Hilbert space. Hereafter, an asterisk denotes a Hermitian conjugate, while a dagger denotes an adjoint.

1. $C^{\pm}$ with a negative metric. We introduce scalar intermediate bosons $C^{\pm}$ with a negative metric in addition to the usual weak vector bosons $B^{\pm}$. The free Lagrangian density of these boson fields is given by

$$\mathcal{L}_0=-\tfrac{1}{2}(\partial_\mu B_\nu{}^\dagger-\partial_\nu B_\mu{}^\dagger)(\partial_\mu B_\nu-\partial_\nu B_\mu)-m_B{}^2B_\mu{}^\dagger B_\mu +\partial_\mu C^\dagger\partial_\mu C+m_B{}^2C^\dagger C, \qquad (2.1)$$

where

$$B_\mu{}^\dagger=(B_k{}^*, -B_4{}^*), \quad C^\dagger=\eta^{-1}C^*\eta,$$
$$[\eta, B_\mu]=0, \qquad \{\eta, C\}=0. \qquad (2.2)$$

For simplicity, we have taken the mass of C, m_C, to be equal to that of B, m_B, although a variation of the C mass keeps most of the following results unchanged (Sec. VII). A fundamental Lagrangian density for the weak interaction is

$$\mathcal{L}_{\rm wk}=gJ_\lambda{}^\dagger[B_\lambda+(i/m_B)\partial_\lambda C]+\text{ adjoint}, \qquad (2.3)$$

where J_λ is the weak current [e.g., $i\bar\psi_e\gamma_\lambda(1+\gamma_5)\psi_{\nu_e}+i\bar\psi_\mu\gamma_\lambda(1+\gamma_5)\psi_{\nu_\mu}$ for leptons] and g is a coupling constant. It is easy to see that the Lagrangian (2.3) implies an effective Fermi coupling of the form (1.1).

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[1] H. Terazawa, Phys. Rev. Letters **22**, 254 (1969); **22**, 442(E) (1969).

2. Haller-Landovitz-Goldberg (HLG) formalism. Haller, Landovitz, and Goldberg[2] have proposed a model of the weak interaction in which the boson field is described as an admixture of spin-1 and spin-0 components. The spin-0 component has a negative metric. The free Lagrangian density in their model is

$$\mathcal{L}_0=-\partial_\mu B_\nu{}^\dagger\partial_\mu B_\nu-m_B{}^2B_\mu{}^\dagger B_\mu, \tag{2.4}$$

where

$$B_\mu{}^\dagger=(B_k{}^*,\,-\eta^{-1}B_4{}^*\eta),\qquad [\eta,B_k]=0,\qquad \{\eta,B_4\}=0, \tag{2.5}$$

while the weak interaction is expressed by

$$\mathcal{L}_{\rm wk}=gJ_\lambda{}^\dagger B_\lambda+\text{ adjoint}. \tag{2.6}$$

3. $\xi=1$ formalism. Lee and Yang[3] have proposed the ξ-limiting formalism in which they have introduced the artifice of a negative metric which makes the parameter ξ take on the role of a regulator. If, for simplicity, we consider the case of $\xi=1$ corresponding to the case of $m_B=m_C$ in Sec. II 1, we obtain the $\xi=1$ formalism which leads to the same effective Fermi interaction (1.1) as the two other formalisms 1 and 2. The free and interaction Lagrangian densities in this model are, respectively,

$$\mathcal{L}_0=-\tfrac{1}{2}(\partial_\mu B_\nu{}^\dagger-\partial_\nu B_\mu{}^\dagger)(\partial_\mu B_\nu-\partial_\nu B_\mu)-(\partial_\mu B_\mu)^\dagger(\partial_\nu B_\nu)-m_B{}^2B_\mu{}^\dagger B_\mu, \tag{2.7}$$

where

$$B_\mu{}^\dagger=(B_k{}^*,\,-\eta^{-1}B_4{}^*\eta),\qquad [\eta,B_k]=0,\qquad \{\eta,B_4\}=0, \tag{2.8}$$

and

$$\mathcal{L}_{\rm wk}=gJ_\lambda{}^\dagger B_\lambda+\text{ adjoint}. \tag{2.9}$$

Although at present it seems impossible to overcome the field-theoretical difficulties of the negative metric,[4] which appears in every state containing a scalar boson or the scalar part of a vector boson, we must at least show that the breakdown of probability conservation can never be seen in any process. Recently Lee and Wick[4] have shown that unstable negative-metric states are consistent with unitarity of the S matrix provided that the Hamiltonian H is self-adjoint, i.e., $H^\dagger=H$. We will follow their interpretation of the negative-metric state.

III. FINITE SELF-MASSES

Our first assumption is that the weak interaction of leptons should be described by one of the three formalisms presented in Sec. II. From this assumption, the weak self-masses of leptons have only logarithmic divergences and are calculated to the lowest order to be

$$\delta m_l{}^{\rm wk}=-(g^2/8\pi^2)m_l\ln(\Lambda^2/m_B{}^2)+\cdots\qquad \text{for } l=e \text{ or } \mu \tag{3.1}$$

and

$$\delta m_\nu{}^{\rm wk}=0\qquad \text{for the neutrinos,} \tag{3.2}$$

where Λ is a cutoff energy. Moreover, we assume[5] that these weak divergences should cancel those of the electromagnetic self-masses of the leptons:

$$\delta m_l{}^{\rm em}=(3e^2/16\pi^2)m_l\ln(\Lambda^2/m_l{}^2)+\cdots, \tag{3.3}$$

so that

$$\delta m_l=\delta m_l{}^{\rm em}+\delta m_l{}^{\rm wk} \tag{3.4}$$

is finite. Then, the coupling constant g can be predicted to be

$$g^2=\tfrac{3}{2}e^2. \tag{3.5}$$

From this relation and the familiar relation

$$g^2/m_B{}^2=G/\sqrt{2} \tag{3.6}$$

between g and the Fermi coupling constant G ($Gm_p{}^2=1.026\times10^{-5}$), the mass of the weak boson can also be determined:

$$m_B=(3\sqrt{2}e^2/2G)^{1/2}\simeq137.7m_p\quad(\simeq m_p/\alpha). \tag{3.7}$$

This mass is so large that a direct test of this model $p+\bar{p}\to B^++B^-$, etc., will not be possible until a colliding-beam machine with beam energies greater than 100 GeV becomes available. From Eqs. (3.1), (3.3)–(3.5), and (3.7), the finite self-masses of the electron and muon become

$$\delta m_l=(3e^2/16\pi^2)m_l\ln(m_B{}^2/m_l{}^2)\qquad(\simeq0.044m_e \text{ for } e), \tag{3.8}$$

which shows that the mass of the weak boson plays the role of an effective cutoff energy in quantum electrodynamics.

We note that this model suggests the following possibility: The weak interaction is not weak but is strong enough to compete with the electromagnetic interaction; that is, the weak interaction behaves like a weak interaction only because the masses of the intermediate bosons are much heavier than those of known particles. We also note that the masses of "weak particles" (leptons and intermediate bosons) and hadrons keep their balance, e.g., $m_Bm_e\simeq m_\pi m_p$.

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[2] K. Haller, L. F. Landovitz, and I. Goldberg, Nuovo Cimento **48**, 303 (1967).

[3] T. D. Lee and C. N. Yang, Phys. Rev. **128**, 885 (1962).

[4] See Ref. 2 and T. D. Lee and G. C. Wick, Nucl. Phys. **B9**, 209 (1969).

[5] A similar idea was proposed about thirty years ago. See E. C. G. Stueckelberg, Nature **144**, 118 (1939); S. Sakata and O. Hara, Progr. Theoret. Phys. (Kyoto) **2**, 30 (1947).

IV. RENORMALIZATION

Since the divergences associated with the weak interaction of leptons proposed in Sec. II are only logarithmic, the theory is renormalizable. In the usual way we introduce the electromagnetic interactions of leptons and weak bosons according to the minimality principle. Charge conservation and gauge invariance hold, as does the Ward identity. All that we need to do is to show that the renormalization of both the electromagnetic and weak interactions of leptons can be consistently accomplished in spite of the existence of the parity-violating interaction.

First, we consider the renormalization of the masses, the wave functions, and the weak coupling constant g for leptons due to the weak interaction. We follow Ioffe's method[6] of renormalization. Let the unrenormalized mass and Green's function of the lepton be m_0 and $G_0(i\gamma p)$, respectively. Then, the equation for the Green's function has the form

$$[i\gamma p+m_0+M(i\gamma p)]G_0(i\gamma p)=1, \tag{4.1}$$

where $\gamma p\equiv\gamma_\mu p^\mu$ and where $M(i\gamma p)$ is the mass operator for which the general expression should have, because of T invariance, the form

$$M(i\gamma p)=i\gamma p M_1(p^2)+i\gamma p\gamma_5 M_2(p^2)+M_3(p^2). \tag{4.2}$$

It is seen from Eq. (4.1) that for $p^2\to -m^2$ (where m is the physical mass), the equation for the Green's function becomes

$$Z_l^{-1}[i\gamma p(1+\lambda\gamma_5)+m']G_0(i\gamma p)\,|_{i\gamma p=-m}=1, \tag{4.3}$$

where

$$Z_l^{-1}=1+M_1(-m^2), \tag{4.4}$$

$$m'=Z_l[m_0+M_3(-m^2)], \tag{4.5}$$

and

$$\lambda=Z_l M_2(-m^2). \tag{4.6}$$

We now make a combination of the transformation and the numerical renormalization of the wave function ψ_0:

$$\psi_0=Z_l^{1/2}S\psi_R, \tag{4.7}$$

where

$$S=A+B\gamma_5,\qquad A=\tfrac{1}{2}[(1+\lambda)^{-1/2}+(1-\lambda)^{-1/2}], \tag{4.8}$$

and

$$B=\tfrac{1}{2}[(1+\lambda)^{-1/2}-(1-\lambda)^{-1/2}]. \tag{4.9}$$

Then, Eq. (4.1) can be transformed into the equation for the renormalized Green's function G_R defined by

$$G_0=Z_l S G_R \tilde{S}, \tag{4.10}$$

where

$$\tilde{S}=\gamma_4 S^*\gamma_4. \tag{4.11}$$

The equation for G_R is

$$\{i\gamma p+m+i\gamma p(1-\lambda\gamma_5)Z_l[M_1(p^2)-M_1(-m^2)]/(1-\lambda^2)$$
$$+i\gamma p(\gamma_5-\lambda)Z_l[M_2(p^2)-M_2(-m^2)]/(1-\lambda^2)$$
$$+Z_l[M_3(p^2)-M_3(-m^2)]/(1-\lambda^2)^{1/2}\}G_R(i\gamma p)=1, \tag{4.12}$$

while the physical mass m is given by

$$m=m'/(1-\lambda^2)^{1/2}. \tag{4.13}$$

For $p^2\to -m^2$, the renormalized Green's function is $G_R\to(i\gamma p+m)^{-1}$, i.e., it has the form of the Green's function of a free particle with mass m. Next, for the purpose of the vertex renormalization, we consider the following expression for the weak-interaction energy, neglecting the other finite terms, e.g., $\sigma_{\mu\nu}q_\nu B_\mu$:

$$ig_0\bar{\psi}_\nu(p'^2)\gamma_\mu[a(p^2,p'^2,k^2)+b(p^2,p'^2,k^2)\gamma_5]\psi_l(p^2)B_\mu(k^2). \tag{4.14}$$

Then, in accordance with the definition of the physical coupling constant g, we should have

$$ig_0\bar{\psi}_\nu(0)\gamma_\mu[a(-m_l^2,0,-m_B^2)$$
$$+b(-m_l^2,0,-m_B^2)\gamma_5]\psi_l(-m_l^2)B_\mu(-m_B^2)$$
$$=ig\bar{\psi}_{\nu R}(0)\gamma_\mu(\alpha+\beta\gamma_5)\psi_{lR}(-m_l^2)B_{R\mu}(-m_B^2) \tag{4.15}$$

and

$$\alpha^2+\beta^2=2, \tag{4.16}$$

where $\psi_{\nu R}$, ψ_{lR}, and $B_{R\mu}$ are the renormalized wave functions of the neutrino, the lepton, and the weak boson, which are related to the unrenormalized ones in Eq. (4.7) and in the following relation:

$$B_\mu=Z_B^{1/2}B_{R\mu}. \tag{4.17}$$

It is worth noticing that there are two different types of solutions for which the ratio of the renormalized parity-conserving interaction to the renormalized parity-nonconserving interaction is equal to the ratio of the corresponding unrenormalized interactions, i.e.,

$$a(-m_l^2,0,-m_B^2):b(-m_l^2,0,-m_B^2)=\alpha:\beta.$$

The solutions are

$$a(-m_l^2,0,-m_B^2)=\pm b(-m_l^2,0,-m_B^2) \tag{4.18}$$

or

$$A_\nu B_l+B_\nu A_l=0. \tag{4.19}$$

In our case of the two-component theory, it can easily be proved[6] that

$$a=b \tag{4.20}$$

and

$$M_1=M_2, \tag{4.21}$$

[6] B. L. Ioffe, Zh. Eksperim. i Teor. Fiz. **38**, 1263 (1960) [Soviet Phys. JETP **11**, 911 (1960)].

so that Eq. (4.18) holds and

$$\lambda = 1 - Z_l. \quad (4.22)$$

From Eqs. (4.7)–(4.9), (4.16), and (4.20)–(4.22), we rewrite Eq. (4.15) as follows:

$$\alpha = \beta = 1 \quad (4.23)$$

and

$$g = g_0 Z_{\rm wk}^{-1} Z_\nu^{1/2} Z_l^{1/2} Z_B^{1/2} (2-Z_\nu)^{-1/2} (2-Z_l)^{-1/2}, \quad (4.24)$$

where the renormalization constant of the weak vertex $Z_{\rm wk}$ is defined by

$$Z_{\rm wk}^{-1} \equiv a = b. \quad (4.25)$$

The renormalization that we have carried out will be valid only in the case in which the quantity λ is less than unity in absolute value, i.e., $|\lambda| \leq 1$. We can easily show, in the same manner as that of Ioffe,[6] that this is actually the case.

Second, we consider the renormalization of the electromagnetic coupling constant e of leptons due to the weak interaction. We repeat the previous procedure. From the expression for the electromagnetic interaction energy (neglecting the other finite terms, e.g., $\sigma_{\mu\nu} q_\nu A_\mu$)

$$ie_0 \bar{\psi}_l(p'^2)\gamma_\mu [c(p^2, p'^2, k^2) + d(p^2, p'^2, k^2)\gamma_5] \psi_l(p^2) A_\mu(k^2), \quad (4.26)$$

we should have, as a condition for consistently renormalizing our model,

$$ie_0 \bar{\psi}_l(-m_l^2)\gamma_\mu [c(-m_l^2, -m_l^2, 0) + d(-m_l^2, -m_l^2, 0)\gamma_5] \psi_l(-m_l^2) A_\mu(0) = ie\bar{\psi}_{lR}(-m_l^2)\gamma_\mu \psi_{lR}(-m_l^2) A_{R\mu}(0), \quad (4.27)$$

where $A_{R\mu}$ is the renormalized wave function of the photon defined by

$$A_\mu = Z_A^{1/2} A_{R\mu}. \quad (4.28)$$

Using Eqs. (4.7)–(4.9) and (4.28), we can transform the condition (4.27) into the relations

$$d/c = \lambda \quad (4.29)$$

and

$$e = e_0 Z_{\rm em}^{-1} Z_l Z_A^{1/2}, \quad (4.30)$$

where the renormalization constant of the electromagnetic vertex $Z_{\rm em}$ is defined by

$$Z_{\rm em}^{-1} \equiv c. \quad (4.31)$$

Now the Ward identity

$$\Gamma_\mu(p^2, p^2, 0) = -i(\partial/\partial p_\mu) G^{-1}(p^2) \quad (4.32)$$

[where $\Gamma_\mu(p^2, p'^2, k^2)$ is the electromagnetic vertex] still holds, and, from Eqs. (4.1), (4.2), and (4.26), it leads to

$$c(p^2, p^2, 0) = 1 + M_1(p^2) \qquad (\text{i.e., } Z_{\rm em} = Z_l) \quad (4.33)$$

and

$$d(p^2, p^2, 0) = M_2(p^2). \quad (4.34)$$

Therefore, in view of Eqs. (4.4) and (4.6), the condition (4.29) is certainly satisfied. Thus, we have accomplished the purpose of this section.

V. ANOMALOUS MAGNETIC MOMENTS OF LEPTONS

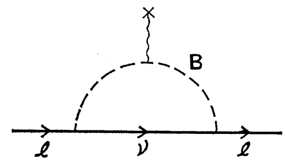

FIG. 1. Diagram for contributions of the weak boson to the anomalous magnetic moments of leptons.

One of several difficulties in the usual weak-boson theory is that physical quantities such as the contribution to the anomalous magnetic moment of leptons are found to be logarithmically divergent. In our model, however, it is finite and is calculated to the lowest order (Fig. 1) to be

$$\Delta(g_l - 2)/2 = -Gm_l^2/12\sqrt{2}\pi^2 \simeq -7.7\times 10^{-10} \qquad (\text{for } \mu). \quad (5.1)$$

Here, for simplicity, we have assumed only the minimal electromagnetic interaction introduced in the HLG formalism stated in Sec. II. This result is one order of magnitude smaller than those previously given by other authors.[7]

VI. WEAK REACTIONS

The simplest weak reaction is a decay of the weak boson. The decay widths for two-lepton modes are easily calculated to be

$$\Gamma_B(B \to l\nu) = Gm_B^3/6\sqrt{2}\pi \simeq 1.0 \text{ GeV}. \quad (6.1)$$

Since the mass m_B ($\simeq 140 m_p$) is so large, the weak boson has so many decay modes, especially hadronic ones whose widths cannot be estimated easily, that the total decay width of the boson may be much larger than $\Gamma_B(B\to l\nu)$, i.e.,

$$\Gamma_B^{\rm tot} \gg 1 \text{ GeV}. \quad (6.2)$$

We therefore expect that this newcomer to particle physics would be an absurdly unstable particle.

[7] N. Byers and F. Zachariasen, Nuovo Cimento **18**, 1289 (1960); R. D. Amado and L. Holloway, *ibid.* **30**, 1083 (1963); **30**, 1572 (1963); G. Segrè, Phys. Letters **7**, 357 (1963); H. Pietechmann, Z. Physik **170**, 409 (1964); R. A. Shaffer, Phys. Rev. **135**, B187 (1964); S. J. Brodsky and J. D. Sullivan, *ibid.* **156**, 1644 (1967).

Next, we consider lepton-lepton scattering. It has been known for a long time[8] that the usual local Fermi interaction leads to lepton-lepton cross sections which increase quadratically and violate the unitarity limit at an energy of about 300 GeV in the center-of-mass system. In our model, however, the interaction being effectively nonlocal, the scattering cross sections are better behaved at high energies. There are two different types of lepton-lepton scattering. One is mediated by the weak boson in the t channel, and the other is mediated by it in the s channel. An example of the former kind is the reaction

$$\nu_\mu+e^-\rightarrow\mu^-+\nu_e. \tag{6.3}$$

Since it consists of all the partial-wave scatterings caused by the propagation of spin-1 or spin-0 particles in the t channel, its cross section asymptotically becomes constant at high energies as follows:

$$\sigma(\nu_\mu+e^-\rightarrow\mu^-+\nu_e)\simeq(G^2m_B{}^2/\pi)\,(1+m_B{}^2/W^2)^{-1} \tag{6.4}$$

$$\underset{W\to\infty}{\longrightarrow}(G^2m_B{}^2/\pi)\ (\simeq 2.8\times10^{-34}\ \text{cm}^2), \tag{6.5}$$

where W is the total energy in the center-of-mass system.

Another example of t-channel exchange, $e^-+\mu^+\rightarrow\nu_e+\bar{\nu}_\mu$, has a similar cross section:

$$\sigma(e^-+\mu^+\rightarrow\nu_e+\bar{\nu}_\mu)\simeq(G^2m_B{}^2/2\pi)\,(1+m_B{}^2/W^2)^{-1}, \tag{6.6}$$

while an example of s-channel exchange, $\bar{\nu}_e+e^-\rightarrow\bar{\nu}_\mu+\mu^-$, has a different high-energy behavior:

$$\sigma(\bar{\nu}_e+e^-\rightarrow\bar{\nu}_\mu+\mu^-)$$

$$\simeq(G^2m_B{}^2/3\pi)\,(m_B{}^2/W^2)\,|\,1-(m_B{}^2/W^2)\,|^{-2} \tag{6.7}$$

$$\underset{W\to\infty}{\longrightarrow}0. \tag{6.8}$$

One of the best tests of this model would be to examine the high-energy behavior of these reactions.

VII. RESULTS AND DISCUSSION

In Sec. III, we have predicted the coupling constant g and the mass m_B to be $g^2=\frac{3}{2}e^2$ and $m_B=137.7m_p$, respectively. A remaining question would be whether these predicted values would be changed by generalizations of the model or by higher-order corrections to the self-masses of leptons. As far as the former problem is concerned, Fukuda[9,10] has generalized our model for an arbitrary value of ξ and obtained the following result: In order to cancel the logarithmic divergences of the self-masses of leptons from electromagnetic and weak interactions, the physical mass of the usual weak boson m_B is arbitrary, but the additional scalar boson with a negative metric has the universal mass proposed by us, i.e., $m_C=137.7m_p$ independently of the value ξ. The latter problem of higher-order corrections is very complex and difficult and has not been solved. It is expected that the relation to determine g^2 will be a quadratic equation for the fourth-order self-masses of leptons. Whether the equation has a proper solution is an important question.

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[8] For example, T. D. Lee and C. S. Wu, Ann. Rev. Nucl. Sci. **15**, 381 (1965).

[9] R. Fukuda (private communication).

[10] Pestieau and Roy considered a similar cancellation within a unified model of leptons and predicted values different from ours for the weak boson masses. See J. Pestieau and P. Roy, Phys. Rev. Letters **23**, 349 (1969).