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Anomalous magnetic moment of electron and constraint on composite weak-interaction bosons

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We examine the possibility of the existence of composite weak-interaction bosons. Constraints from the anomalous magnetic moment of the electron are used to determine upper limits on the couplings of the bosons, an entire spectrum of which may exist in the composite model. It is found that the coupling of a singlet or triplet pseudoscalar to leptons is very small. This explains the V,A nature of weak interactions. It is assumed that composite bosons have a structure similar to mesons where the constituents are preons instead of quarks. The present experimental status, although in perfect agreement with the standard model, does not rule out composite bosons. We consider the possible signals for compositeness, including the anomalous $e^+e^-\gamma$ events.

I. INTRODUCTION

The observation[1] of the $W^{\pm}$ and the Z^0 bosons is a clear verification of the $SU_L(2)\times U(1)$ theory of electroweak interactions. There are some aspects of the theory, however, there are not so satisfactory. The question that still remains is the existence of elementary scalars (as in the standard model[2]) called Higgs bosons. They are necessary in order to generate the mass of the gauge bosons. These scalars have a new fundamental interaction, and couple to fermions and gauge bosons. Within the standard model, there are no constraints on the self-coupling of the Higgs bosons or on its Yukawa coupling to fermions. There is also the so-called "naturalness problem"[3] associated with the existence of elementary scalars. Though supersymmetry does solve the naturalness problem, the lack of constraints on the Higgs-boson coupling still remains.

The symmetry $SU_L(2)\times U(1)$ can however also be dynamically broken. In such a scheme, called the "technicolor model"[4] the Higgs particles are scalar composites of a new set of fermions, bound by a new non-Abelian gauge interaction. This assumption of composite scalars solves both the problems of determining the coupling and naturalness. This interaction is required to be very strong and is assumed to be of the QCD type with a scale of a TeV or so.

It has also been speculated[5] that the weak-interaction boson are composite. In such models, spontaneous breaking of $SU(2)\times U(1)$ symmetry is not needed to generate the mass of the bosons. The naturalness problem is thus automatically avoided. Several attempts have been made to develop such a model that regards weak interactions as essentially residual. The compositeness scale could be as low as a few hundred GeV to a few TeV or more. If the composite scale is small (~ 100 GeV), the mass of the bosons would be consistent with the standard model. However, if the scale of compositeness is higher, understanding the mass of the bosons would be a problem. In view of the verification[1] of the $W^{\pm}$ mass, we assume the mass of the vector boson is indeed 83 GeV.

There appears to be no clear indication as yet of the presence of any compositeness of these bosons. An important consequence of such compositeness would be the spectrum of bosons of various spins. Note that in some models[6] that obey the "complementary principle," there are restrictions on the spectrum of bosons that appear. The purpose of this paper is to examine the possibility of the existence of this spectrum and the technique of observing it. In order to observe these bosons one should know their coupling to fermions and coupling among themselves. The anomalous magnetic moment of the electron provides a good constraint on these couplings. In Sec. II of the paper, we derive upper limits on the couplings from restrictions on the anomalous magnetic moment. Note that these limits would be independent of the model of compositeness. We shall consider only a spectrum of bosons: scalar, pseudoscalar, vector, and axial vector. Our calculations, however, are not for a gauge theory. Indeed, there is no satisfactory (gauge) theory of composite bosons. We therefore consider it safer to use effective local interactions.

We emphasize two essential requirements that any composite model should meet: the $V-A$ structure of weak interactions and the far more well-established requirement of $SU_L(2)\times U(1)$ symmetry. These requirements are in addition to any other requirement[7] that an underlying composite model may be expected to satisfy. Further, there are also constraints on composite models. Most of these constraints are model dependent (see Refs. 7 and 8) and are not considered here. All we point out is that there are model-independent (but depend on parametrization), highly restrictive[9] bounds on composite leptons from $e^+e^-\to e^+e^-$, $\mu^+\mu^-$, $\nu\bar{\nu}$. However, there is no good restriction on the compositeness of weak bosons. Measurements of charge asymmetry are in agreement with the standard model. These should not be altered by the presence of pseudoscalar and scalar bosons. The signature of the excited states of bosons, however, would alter the forward-backward asymmetry from the standard model at energies close to the excited states. It may also be expected that the composite bosons may show up in neutrino counting[10] near the Z^0 peak. This signature, also, may not be large enough, as we shall demonstrate by examining the restrictions from $(g-2)$ of the electron.

Before ending this section, we would like to point out that most of the existing models[8] of compositeness would not survive if only a triplet of bosons exists. However, if we require that the preons, constituents of the weak bosons, are bound by SU(2)-color interaction and that there is no degeneracy of the lightest preon (denoted by α), we have only three approximately degenerate lightest states: $\alpha\bar{\alpha}$, $\alpha\alpha$, $\bar{\alpha}\bar{\alpha}$. We also point out that no subcomponent model to date is able to explain consistently all problems faced by the standard model. Clearly, more experimental input is required to understand the real composite structure—if any exists. The observation of the (pseudo)scalar state close to W and Z is a necessary requirement.

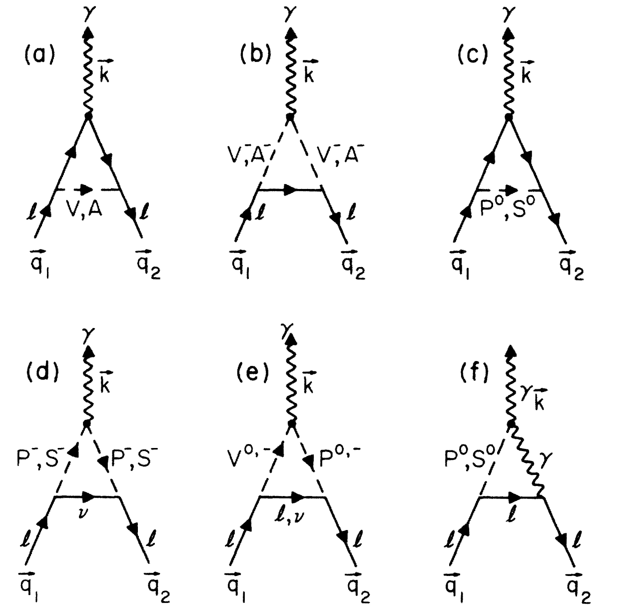

FIG. 1. One-loop graphs that contribute to the anomalous magnetic moment of a lepton.

II. CONTRIBUTION TO THE ANOMALOUS MAGNETIC MOMENT OF LEPTON

In this section we evaluate the correction to the anomalous magnetic moment of the lepton (a_l) arising from the presence of bosons that couple to leptons. We assume that there exists a spectrum of bosons: scalars (S), pseudoscalars (P), vector (V), and axial vector (A). The one-loop graphs that can contribute to the anomalous magnetic moment arising from such bosons are shown in Fig. 1.

The contribution from these graphs can be written as

$$\bar{u}(q_2)e\Gamma_\mu(q_2,q_1)u(q_1)=\bar{u}(q_2)\left[e\gamma_\mu F_1(k^2)+\frac{ie}{2m}\sigma_{\mu\nu}k^\nu F_2(k^2)+(\gamma_5 \text{ terms})\right]u(q_1)\ . \tag{1}$$

The anomalous magnetic moment is defined as $F_2(k=0)$. Since our model of composite bosons is not a gauge theory, $F_2(k^2)$ need not be finite. It is for this reason that we exclude contributions from tensor bosons, if any exist. In order to evaluate contributions to the anomalous magnetic moment we assume effective local interactions for the gauge bosons. However, the use of effective local interactions gives divergent results. Thus, we regularize by introducing a cutoff Λ. We follow the standard procedure of replacing the propagator

$$\frac{1}{p^2+i\epsilon}\to-\int_0^{\Lambda^2}\frac{d\lambda}{(p^2-\lambda+i\epsilon)^2} \tag{2}$$

in order to introduce the cutoff. The resulting expressions are finite for finite Λ, but generally infinite and meaningless for $\Lambda\to\infty$. The presence of this cutoff Λ, in our results, can be understood if Λ is regarded as a composite scale. For a momentum larger than Λ, the effective couplings for bosons are no longer valid. These effective couplings then have to be replaced by interactions of subcomponents. An analogous situation is found in hadron interactions, where effective local interactions adequately describe low-energy phenomena, but at energies above the hadron composite scale it is necessary to consider the interactions of quarks.

We use m to denote the mass of the lepton, M to denote the mass of the boson, the subscripts of M to denote the spin-parity, and Λ to denote the composite scale. The Lagrangians for the local interactions are listed in Table I. In the technicolor model, the pseudoscalar Higgs boson couples to the lepton with a derivative coupling. We, therefore, consider both cases, with and without derivative coupling, in evaluating the anomalous magnetic moment when a pseudoscalar contributes. By looking at the interaction Hamiltonians, one can discern the dimensions of the couplings. In particular $g_{Sll'}$, $g_{Pll'}$, $g_{Vll'}$, and $g_{All'}$ are dimensionless, whereas $g'_{Pll'}$, $g_{PV\gamma}$, $g_{P\gamma\gamma}$, and $g_{S\gamma\gamma}$ have a dimension of mass inverse. Figures 1(a)—1(d) have been

TABLE I. Effective local interaction for gauge bosons.

$H_{ll'S}=g_{Sll'}\bar{\psi}_l\psi_l S$
$H_{ll'P}=g_{Pll'}\bar{\psi}_l\gamma_5\psi_{l'}P$ or
$\quad=g'_{Pll'}\bar{\psi}_l\gamma_\mu\gamma_5\psi_{l'}(\partial^\mu-ieA^\mu)P$
$H_{ll'V}=g_{Vll'}\bar{\psi}_l\gamma_\mu\psi_{l'}V^\mu$
$H_{ll'A}=g_{All'}\bar{\psi}_l\gamma_\mu\gamma_5\psi_{l'}A^\mu$
$H_{VP\gamma}=\sqrt{4\pi\alpha}g_{PV\gamma}P\,{}^*F^{\alpha\beta}(\partial_\alpha-ieA_\alpha)V_\beta$
$H_{P\gamma\gamma}=g_{P\gamma\gamma}\,{}^*F^{\mu\nu}F_{\mu\nu}P$
$H_{S\gamma\gamma}=2g_{S\gamma\gamma}F^{\mu\nu}F_{\mu\nu}S$

evaluated earlier for the limit $\Lambda\to\infty$ (where such a limit is meaningful) with a nonderivative coupling, in Ref. 11 and references cited therein. We have repeated these calculations for arbitrary Λ and our results agree with them for $\Lambda\to\infty$. We have evaluated the graphs of Figs. 1(c) and 1(d) with a derivative coupling for pseudoscalar. Graphs of Figs. 1(e) and 1(f) have already been used to determine the restriction[12] on compositeness for large values of Λ ($\Lambda=1$ TeV). We would like to point out that there will be no contribution from the diagram analogous to Figs. 1(e) and 1(f) with axial-vector and scalar coupling, instead of vector and pseudoscalar, due to charge conjugation. Within the standard model, therefore, there is no contribution by the graph with Z^0-Higgs-photon vertex.

We assume that the difference between the experimental and theoretical values of the anomalous magnetic moment[13] of the electron ($\sim 2\times10^{-10}$) is saturated by the contribution from Figs. 1(a)—1(f). The constraints on coupling, thus evaluated, are very approximate upper limits and expected values would be much lower. It must be noted that there are several contributions to the anomalous magnetic moment. In particular, there should be contributions[14] from composite leptons of $O(m/\Lambda)$.

The contribution to a_l from Fig. 1(a) is

$$\Delta a_l=\frac{m^2}{4\pi^2}\int_0^1 dy\left\{\frac{g_{Vll}{}^2y^2(1-y)}{m^2y^2+M_V{}^2(1-y)}+g_{All}{}^2\left[\frac{3}{M_A{}^2}(y-\tfrac{3}{2}y^2)\ln\left[\frac{m^2y^2+(M_A{}^2+\Lambda^2)(1-y)}{m^2y^2+M_A{}^2(1-y)}\right]\right.\right.$$
$$\left.\left.+\frac{\Lambda^2y(1-y)[2m^2y^3+(1-y)(4-y)M_A{}^2]}{M_A{}^2[m^2y^2+(M_A{}^2+\Lambda^2)(1-y)][m^2y^2+M_A{}^2(1-y)]}\right]\right\}. \qquad (3)$$

It is interesting to note that the term with the logarithm is finite and becomes zero if $m^2\ll M_A{}^2,\Lambda^2$ (as can be seen by integrating over y). In such a limit we have

$$\Delta a_l=\frac{m^2}{4\pi^2}\int_0^1 dy\left[\frac{g_{Vll}{}^2y^2(1-y)}{m^2y^2+M_V{}^2(1-y)}+\frac{\Lambda^2g_{All}{}^2y(1-y)(4-y)}{(M_A{}^2+\Lambda^2)[m^2y^2+M_A{}^2(1-y)]}\right]. \qquad (4)$$

In the limit $m\ll M_{V,A}$ we can evaluate the integrals exactly to get

$$\Delta a_l=\frac{m^2}{12\pi^2}\left[\frac{g_{Vll}{}^2}{M_V{}^2}+\frac{5g_{All}{}^2}{2M_A{}^2}\frac{\Lambda^2}{M_A{}^2+\Lambda^2}\right]. \qquad (5)$$

The constraints on the couplings g_{Vee} and g_{Aee} are shown in Fig. 2.

The contributions from both V^- and A^- in Fig. 1(b) to the anomalous magnetic moment are equal and are given by

$$\Delta a_l=\frac{m^2}{8\pi^2}\int_0^1 dy\left[\frac{g_{Vl\nu}{}^2\Lambda^2y^2(1-y)[2(1+y)-(m/M_V)^2(1-y)]}{[m^2y^2+(M_V{}^2-m^2)y+\Lambda^2(1-y)][m^2y^2+(M_V{}^2-m^2)y]}+V\to A\right]. \qquad (6)$$

The limits on the coupling $g_{Ve\nu}(g_{Ae\nu})$ is shown in Fig. 3. In the limit $m^2\ll M_V{}^2,\Lambda^2$ one has, with $\lambda=(M_V/\Lambda)^2$,

$$\Delta a_l=\frac{m^2}{4\pi^2}\left[\left[\frac{-5(\lambda+1)+4\lambda^2}{6(\lambda-1)^3}+\frac{\lambda(2+\lambda)}{(\lambda-1)^4}\ln\lambda\right]\frac{g_{Vl\nu}{}^2}{M_V{}^2}+V\to A\right]_{\lambda\to0}=\frac{5m^2}{24\pi^2}\left[\frac{g_{Vl\nu}{}^2}{M_V{}^2}+V\to A\right]. \qquad (7)$$

The contribution from Fig. 1(c) which corresponds to the familiar scalar Higgs boson and a pseudoscalar is (for nonderivative coupling)

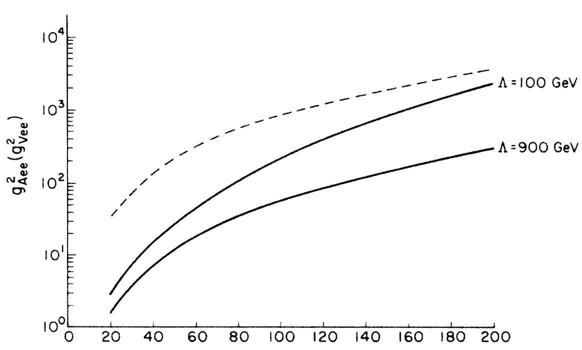

FIG. 2. The constraints on the couplings $g_{Aee}{}^2$ and $g_{Vee}{}^2$ vs the mass of bosons A and V. The allowed values are below the curve. The solid curve is for A (axial vector), whereas the dashed curve is for V (vector).

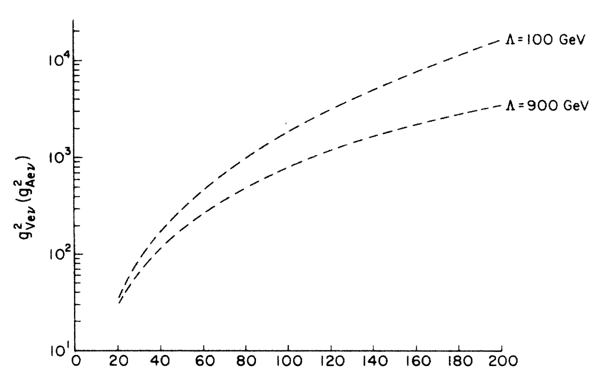

FIG. 3. The constraints on the couplings $g_{Ae\nu}{}^2$ and $g_{Ve\nu}{}^2$ vs the mass of bosons A and V. The allowed values are below the curve.

$$\Delta a_l = \frac{m^2}{8\pi^2}\int_0^1 dy \left[\frac{g_{Sll}{}^2 y^2(2-y)}{m^2y^2+M_S{}^2(1-y)} + \frac{g_{Pll}{}^2 y^3}{m^2y^2+M_P{}^2(1-y)} \right] . \tag{8}$$

If $m^2 \ll M_S{}^2, M_P{}^2$, we have a rather simple relation:

$$\Delta a_l = \frac{m^2}{8\pi^2}\left\{ \frac{g_{Sll}{}^2}{M_S{}^2}\left[\ln\left[\frac{M_S}{m}\right]^2 - \frac{7}{6}\right] + \frac{g_{Pll}{}^2}{M_P{}^2}\left[\ln\left[\frac{M_P}{m}\right]^2 - \frac{11}{6}\right]\right\} . \tag{9}$$

If however we assume a derivative coupling for a pseudoscalar we get

$$\Delta a_l = \frac{m^2}{8\pi^2}\int_0^1 dy \left[\frac{g_{Sll}{}^2 y^2(2-y)}{m^2y^2+M_S{}^2(1-y)} + \frac{g'_{Pll}{}^2 (m/M_P)^2\Lambda^2 y^4(1-y)}{[m^2y^2+(M_P{}^2+\Lambda^2)(1-y)][m^2y^2+M_P{}^2(1-y)]} \right] . \tag{10}$$

The restrictions on the couplings g_{See} and g_{Pee} are shown in Fig. 4. Clearly these restrictions are not very good. Better restrictions[15] on g_{See} are given by $e^+e^- \to \mu^+\mu^-$. For $M_S \sim 50$ GeV one gets $(g_{See})^2 < 0.003$.

Figure 1(d) gives the contributions from P^- and S^- to a_l. With a nonderivative coupling, the contributions are equal, and given by

$$\Delta a_l = \frac{m^2}{8\pi^2}\int_0^1 dy \left[\frac{g_{Sl\nu}{}^2 \Lambda^2 y^2 (1-y)^2}{[m^2y^2+(M_S{}^2-m^2)y+\Lambda^2(1-y)][m^2y^2+(M_S{}^2-m^2)y]} + S \to P \right] . \tag{11}$$

We note that the integral can be evaluated easily in the limit $m^2 \ll M_S{}^2, \Lambda^2$ giving us, with $\lambda = (M_S/\Lambda)^2$,

$$\Delta a_l = \frac{m^2}{8\pi^2}\left[\left[\frac{1-5\lambda-2\lambda^2}{6(\lambda-1)^3} + \frac{\lambda^2}{(\lambda-1)^4}\ln\lambda \right] \frac{g_{Sl\nu}{}^2}{M_S{}^2} + S \to P \right]_{\lambda\to 0} = \frac{m^2}{48\pi^2}\left[\frac{g_{Sl\nu}{}^2}{M_S{}^2} + S \to P \right] . \tag{12}$$

With a derivative coupling, however, we get, for the pseudoscalar,

$$\Delta a_l = \frac{m^2}{8\pi^2}\int_0^1 dy\, g'_{Pl\nu}{}^2 \left[3(2y-y^3-9y^2/2)\ln\left[\frac{m^2y^2+(M_P{}^2-m^2)y+\Lambda^2(1-y)}{m^2y^2+(M_P{}^2-m^2)y}\right] + \frac{m^2\Lambda^2 y^2(1-y)^4}{[m^2y^2+(M_P{}^2-m^2)y+\Lambda^2(1-y)][m^2y^2+(M_P{}^2-m^2)y]} \right] . \tag{13}$$

Figure 5 gives the restrictions on $g_{Pe\nu}$ ($g_{Se\nu}$). Again better restrictions are available[15] on $g_{Se\nu}$ from $e\nu_l \to e\nu_l$.

Figure 1(e) gives a contribution to a_l that is slightly different for charged and neutral modes. With a nonderivative coupling we get

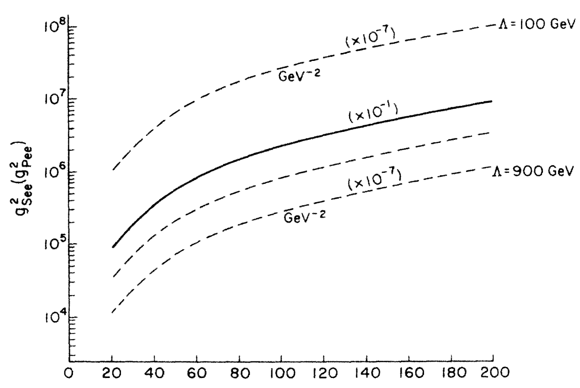

FIG. 4. The constraints on the couplings $g_{See}{}^2$ and $g_{Pee}{}^2$ vs the mass of bosons S and P. The allowed values are below the curve. The solid curve is for a pseudoscalar with a nonderivative coupling and dashed curve for a scalar. The dashed curves with a scale factor $(\times 10^{-7})$ are for a pseudoscalar with a derivative coupling. The couplings are multiplied by the scale factor and then plotted.

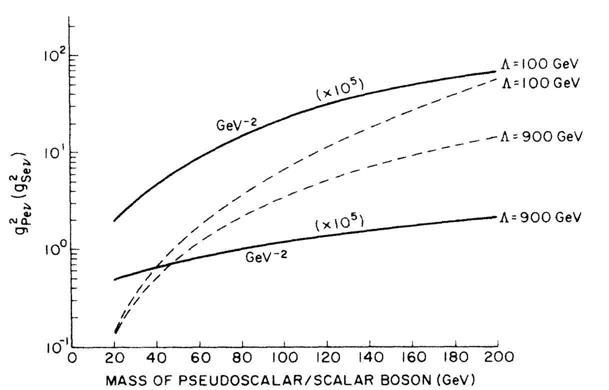

FIG. 5. The constraints on the couplings $g_{Pe\nu}{}^2$ and $g_{Se\nu}{}^2$ vs the mass of bosons S and P. The allowed values are below the curve. The dashed curves give the values with a nonderivative coupling. The solid curves give the values with a derivative coupling. The couplings are multiplied by scale factor and then plotted.

$$\Delta a_l = \frac{g_P g_{PV\gamma} g_V}{8\pi^2} m \int_0^1 dy\, y \int_0^1 dx \left[\ln \left\{ \frac{A + \Lambda^2(1-y)}{A} \right\} - \frac{2m^2(1-y)^3\Lambda^2}{[A+\Lambda^2(1-y)]A} \right] , \tag{14}$$

where

$A = m^2y^2 - 2m^2y + (M_P{}^2 - M_V{}^2)xy + M_V{}^2 y + m^2$ for the neutral mode ,

$A = m^2y^2 - m^2y + (M_P{}^2 - M_V{}^2)xy + M_V{}^2 y$ for the charged mode .

A derivative coupling gives us a slightly different result:

$$\Delta a_l = \frac{g_P' g_{PV\gamma} g_V}{4\pi^2} m^2 \int_0^1 dy\, y(1-y) \int_0^1 dx \left[\ln \left\{ \frac{A + \Lambda^2(1-y)}{A} \right\} - \frac{m^2(1-y)^3\Lambda^2}{[A+\Lambda^2(1-y)]A} \right] . \tag{15}$$

Restrictions on the couplings $g_{Pee} g_{PV\gamma} g_{Vee}$ and $g_{Pe\nu} g_{PV\gamma} g_{Ve\nu}$ are given in Fig. 6.

The contributions to the anomalous magnetic moment a_l from graphs of Fig. 1(f) are equal for nonderivative couplings and are

$$\Delta a_l = \frac{g_S g_{S\gamma\gamma} m}{8\pi^2} \int_0^1 dy\, y^2 \int_0^1 dx \left[\ln \left\{ \frac{A + \Lambda^2(1-y)}{A} \right\} - \frac{2m^2\Lambda^2x^2y^2(1-y)}{[A+\Lambda^2(1-y)]A} + S \to P \right] , \tag{16}$$

where

$A = m^2x^2y^2 + M_S{}^2(1-x)y$

with a derivative coupling, the contribution from pseudoscalar becomes

$$\Delta a_l = \frac{g_P g_{P\gamma\gamma} m^2}{4\pi^2} \int_0^1 dy\, y \int_0^1 dx (1-xy) \left[\ln \left\{ \frac{A + \Lambda^2(1-y)}{A} \right\} - \frac{m^2\Lambda^2x^2y^2(1-y)}{[A+\Lambda^2(1-y)]A} \right] . \tag{17}$$

Restrictions of the couplings $g_{Pee} g_{P\gamma\gamma}$ ($g_{See} g_{S\gamma\gamma}$) are shown in Fig. 7.

Finally, we note that the limit $\Lambda \to \infty$ is meaningful (in our formulation) only for graphs that appear in the standard model and its extensions with charged Higgs bosons. These cases have been considered in detail in Ref. 11 and references therein.

III. HOW TO BEST DETECT COMPOSITENESS?

In this section we examine various possibilities of observing compositeness of bosons. Earlier there had been much interest in the anomalous $l^+l^-\gamma$ events and the monojet events. As no more anomalous events have been reported[16] by the UA1 and UA2 Collaborations, the statistical significance of such events goes down. There have been several papers suggesting[12,17] the composite nature of bosons in view of these events. Here we evaluate the decay rate of a vector and a pseudoscalar boson to $l^+l^-\gamma$. Unlike Refs. 12 and 17 our results are not only model independent, but also include a range of values for Λ and M_P. We consider the decay through an intermediate boson state and denote the decay rate as

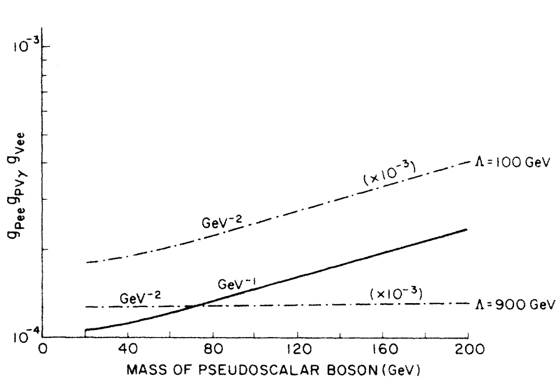

FIG. 6. The constraints on the coupling $g_{Pee} g_{PV\gamma} P_{Vee}$ vs the mass of the pseudoscalar boson. The allowed values are below the curve. The solid curve is with a nonderivative coupling whereas the dot-dashed curve is with a derivative coupling. The couplings are multiplies by the scale factor and then plotted.

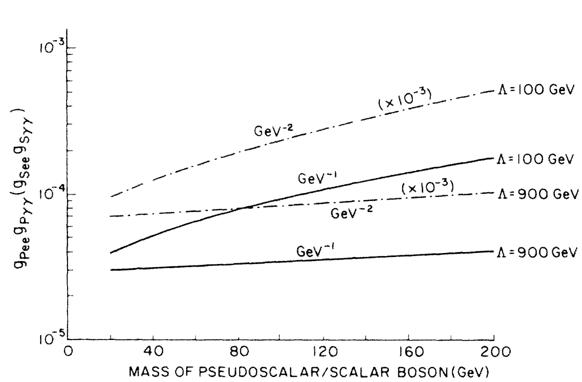

FIG. 7. The constraints on the coupling $g_{Pee} g_{P\gamma\gamma}$ vs the mass of the pseudoscalar boson. The allowed values are below the curve. The solid curve is with a nonderivative coupling whereas the dot-dashed curve is with a derivative coupling. The couplings are multiplied by the scale factor and then plotted.

Γ_I. We have, for a vector,

$$\Gamma_I(V\rightarrow l^+l^-\gamma)=\frac{g_{VP\gamma}{}^2 g_{Pee}{}^2}{12(2\pi)^5 M_V M_P{}^4}\int\frac{d^3k\,d^3p_+d^3p_-}{\omega_k E_+E_-}\delta^4(q-k-p_++p_-)\frac{(q\cdot k)^2(p_+\cdot p_- - m^2)}{1-\dfrac{(p_++p_-)^2}{M_P{}^2}}\,, \tag{18}$$

where we have taken contribution only from an intermediate pseudoscalar. Assuming chiral invariance the contribution from the scalar would be identical. The corresponding relation for a pseudoscalar is

$$\Gamma_I(P\rightarrow l^+l^-\gamma)=\frac{g_{P_sV\gamma}{}^2 g_{Vee}{}^2}{12(2\pi)^5 M_P M_V{}^4}\int\frac{d^3kd^3p_+d^3p_-}{\omega_k E_+E_-}\times\frac{\delta^4(q-k-p_+p_-)}{[1-(p_++p_-)^2/M_V{}^2]^2}[(q\cdot k)^2(m^2-p_+-p_-)$$
$$+2p_+\cdot kp_-\cdot k(m_P{}^2-2q\cdot k)]\,. \tag{19}$$

Using these rates, and also

$$\Gamma(V\rightarrow l^+l^-)=\frac{g_{Vee}{}^2}{6\pi}M_V,\quad \Gamma(P\rightarrow l^+l^-)=\frac{g_{Pee}{}^2}{8\pi}M_P \tag{20}$$

we get the relations

$$\Gamma_I(V\rightarrow l^+l^-)\Gamma(V\rightarrow l^+l^-\gamma)=\frac{g_{VP\gamma}{}^2 g_{Pee}{}^2 g_{Vee}{}^2 M_V A(M_P)}{6\pi(4\pi M_V)^3 M_P{}^4} \tag{21}$$

and

$$\Gamma_I(P\rightarrow l^+l^-)\Gamma(P\rightarrow l^+l^-\gamma)=\frac{g_{VP\gamma}{}^2 g_{Vee}{}^2 g_{Pee}{}^2 M_P A(M_V)}{8\pi(4\pi M_P)^3 M_V{}^4}\,, \tag{22}$$

where $A(M_P)$ and $A(M_V)$ represent integrals over phase space in the expressions for $\Gamma(P\rightarrow l^+l^-\gamma)$ and $\Gamma(V\rightarrow l^+l^-\gamma)$.

In Sec. II we obtained constraints on $g_{VP\gamma}g_{Vee}g_{Pee}$ from $(g-2)$ of the electron. Using these constraints we obtain a limit on the anomalous $l^+l^-\gamma$ events for both a vector and a pseudoscalar particle. In Fig. 8 we plot $R\times\Gamma^2[V/(P)\rightarrow l^+l^-]$ versus the mass of the pseudoscalar, with R defined as

$$R=\frac{\Gamma[V/(P)\rightarrow l^+l^-\gamma]}{\Gamma[V/(P)\rightarrow l^+l^-]}\,. \tag{23}$$

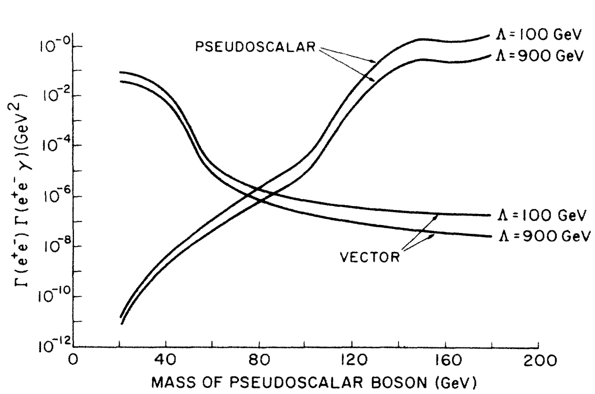

FIG. 8. Constraints on the decay rate $\Gamma(e^+e^-)\Gamma(e^+e^-\gamma)$ vs the mass of the pseudoscalar boson. The expected values are below the curve.

As can be seen the constraints are rather severe for vectors, but not so severe for pseudoscalars if $M_P\geq 90$ GeV, and vice versa for $M_P\leq 90$ GeV. These constraints are model independent. However, contributions from radial excitations are excluded. In that sense, these constraints are approximate. It may be noted[18] that if the subcomponents have charge $(q=)\ \frac{1}{2}$, the radiative decay of the triplet members are forbidden. We denote the triplet of $V^{\pm,0}$ and $P^{\pm,0}$ for vectors and a pseudoscalar, respectively, and the singlet states are denoted by V_s and P_s. However, similar constraints are applicable to anomalous events in $V^-/(P)^-\rightarrow l^-\nu$ in models where subcomponents with $q\neq\frac{1}{2}$, or in models with only three states (see Sec. I) where $V^0\rightarrow P^0\gamma$ and $P^0\rightarrow V^0\gamma$ are forbidden.

Many more interesting constraints can be drawn if we restrict ourselves to a specific composite model. If we assume the dominant contribution to the $PV\gamma$ vertex is a magnetic moment transition, i.e., the constituents have spin $\frac{1}{2}$, then the coupling $g_{PV\gamma}$ can be evaluated in the specific model. For the neutral states $g_{VP\gamma}$ should be $\sqrt{4\pi\alpha}/m_h$, where m_h is the mass of constituents and have charges $\frac{1}{2}$. If the charges are q and $(1-q)$, radiative couplings of charged state are allowed, and $g_{V^-P^-\gamma}$ should be $\sqrt{4\pi\alpha}(2q-1)/m_h$. In a model with only three states and $q=\frac{1}{2}$, $g_{VP\gamma}=0$ but $g_{V^-P^-\gamma}\sim\sqrt{4\pi\alpha}/m_h$. In Sec. II the constraint for a nonderivative coupling was obtained as

$$\frac{g_{V^-e\nu}g_{V^-P^-\gamma}g_{P^-e\nu}}{\sqrt{4\pi\alpha}}\leq 10^{-4}\,. \tag{24}$$

The standard model gives $g_{Vl\nu}{}^2\sim 5.19\times10^{-2}$ which seems to be in agreement with the experiment. Using this value and the above values for $g_{V^-P^-\gamma}$, we get

$$g_{P^-e\nu}\leq 1.8\times10^{-2}$$
$$\text{(model with only three states, } q\neq\tfrac{1}{2})$$
$$\leq\frac{1.8\times10^{-2}}{(2q-1)}$$
$$\text{(models with four states, } q\neq\tfrac{1}{2})\,. \tag{25}$$

The corresponding branching ratios are thus

$$\Gamma(P^{\pm}\rightarrow l^{\pm}\bar{\nu}) \leq 1.35\times 10^{-5} M_P \text{ GeV}$$

(models with only three states)

$$\leq \frac{1.35\times 10^{-5}}{(2q-1)^2} M_P \text{ GeV}$$

$$(\text{model with four states, } q\neq \tfrac{1}{2})\ . \quad (26)$$

These are only of the order of an MeV or less and should be compared with $\Gamma(W^{\pm}\rightarrow l^{\pm}\nu)\sim 230$ MeV. This explains the apparent $V-A$ nature of weak interactions. For a derivative coupling these constraints are no good (a factor of 10^3 larger in couplings).

The coupling $g_{S/P\gamma\gamma}$ should be related[19] to $g_{S/PZ\gamma}$ by the vector-dominance model, i.e.,

$$g_{S/P\gamma\gamma}\sim g_{S/PZ\gamma}\sin\theta_W \sim \sin^2\theta_W \frac{\sqrt{4\pi\alpha}}{m_h}\ . \quad (27)$$

Thus

$$g_{S/P\gamma\gamma}\sim \frac{2\sin^2\theta_W\sqrt{4\pi\alpha}}{M_P} \quad (\text{for singlet states, } q=\tfrac{1}{2})$$

$$\sim \frac{2\sin^2\theta_W\sqrt{4\pi\alpha}}{M_P}(2q-1)$$

$$(\text{for isovector states, } q\neq\tfrac{1}{2})\ . \quad (28)$$

Once again for a nonderivative coupling we obtained, in Sec. II,

$$g_{S/Pee}\, g_{S/P\gamma\gamma} \leq 10^{-4}\ . \quad (29)$$

Using the above constraints we get

$$g_{S_s/P_s ee} \leq \frac{10^{-4} M_P}{2\sin^2\theta_W\sqrt{4\pi\alpha}} \leq 6.6\times 10^{-4} M_P\ ,$$

$$g_{S^0/P^0 ee} \leq 6.6\times 10^{-4}(2q-1)M_P\ . \quad (30)$$

Therefore,

$$\Gamma(P/S\rightarrow l^+l^-) \leq 1.74 M_P{}^3\times 10^{-8} \text{ GeV}$$

$$(\text{for isovector states, } q\neq\tfrac{1}{2})\ . \quad (31)$$

If $M_P\sim 100$ GeV, we have $\Gamma[P/(S)\rightarrow l^+l^-]\leq 10$ MeV as compared to $\Gamma(Z\rightarrow l^+l^-)$ which according to the standard model is ~ 90 MeV.

We thus find that the constraints of $(g-2)$ of a lepton, together with standard subcomponent models, require that pseudoscalar or scalar states may not couple to leptons. This explains the V,A nature of weak interactions. Note that these constraints also apply to the neutral isoscalar. The coupling of fermions to the (pseudo)scalar-triplet states being small, can be explained as a consequence of chiral symmetry and the corresponding PCAC (partial conservation of axial-vector current) relation anologus to the Goldberg-Treiman relation. Similar reasoning for the isoscalar does not hold.

The isovector states couple to quarks and leptons only. For the (pseudo)scalar triplet quark couplings must be dominant. In view of the above argument of chiral-symmetry breaking, the coupling to fermions may be regarded as proportional to the mass of the fermion ($\sim m_f/M_P$). The decay of $S(P)^{\pm}\rightarrow t\bar{b}$, $b\bar{t}$, etc., would be the best signals for its detection. For the $(S/P)^0\rightarrow t\bar{t}, b\bar{b}$ decay rates would be larger. These states would be relatively narrower. The (pseudo)scalar-singlet state would decay dominantly in two gluons.

The decay rate of a singlet pseudoscalar $\Gamma_I(P_s\rightarrow l^+l^-\gamma)$ was evaluated earlier. For anomalous events,

$$\Gamma_I(P_s\rightarrow l^+l^-\gamma) = \frac{g_{PZ\gamma}{}^2 g_{Zee}{}^2 A(M_V)}{12(2\pi)^5 M_P M_V{}^4}\ . \quad (32)$$

We have evaluated $g_{PZ\gamma}$ and g_{Zee} earlier; thus,

$$\Gamma_I(P_s\rightarrow l^+l^-\gamma)\sim \frac{4.3\times 10^{-19}}{M_P} 8\times 10^{16}$$

$$\sim \frac{3.4\times 10^{-2}}{M_P} \text{ GeV}\ . \quad (33)$$

Also for a nonderivative coupling,

$$\Gamma(P_s\rightarrow l^+l^-) \leq 1.7 M_P{}^3\times 10^{-8} \text{ GeV}\ , \quad (34)$$

$$R=\frac{\Gamma_I(P_s\rightarrow l^+l^-\gamma)}{\Gamma(P_s\rightarrow l^+l^-)} \geq \frac{2\times 10^5}{M_P{}^4}$$

$$\geq 0.03 \quad \text{for } M_P = 90 \text{ GeV}\ . \quad (35)$$

Thus the composite model with a nonderivative coupling cannot be ruled out from CERN data. A composite model with a derivative coupling for a pseudoscalar is even harder to rule out. However, in such a model the rate $\Gamma(P_s\rightarrow l^+l^-)$ need not necessarily be small and the V,A structure may not be easily explainable. Combining the results from the CERN UA1 and UA2 Collaborations,[20] one has a total of three anomalous events. Assuming these events are from Z decays, one still has about 4% anomalous events.

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