CERN-TH.4700/87
PM87/19

ONE-LOOP COMPOSITENESS CONTRIBUTIONS TO THE PROPAGATORS OF ELECTROWEAK BOSONS

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A B S T R A C T

We evaluate, by using dimensional regularization, various leading one-loop compositeness contributions to the electroweak boson propagators. We show that the finiteness of some two-point correlators at zero momentum transfer requires gauge-type relations between the three- and four-point effective boson vertices. This finiteness condition also needs some additional effects on loops involving composite particles with higher spins (2, 3/2...). We apply our results to the Z-boson mass shift. We find that an eventual agreement of the measured Z-mass with the standard model predictions within an accuracy of 1%.. feasible at LEP I/SLC provides useful constraints on the size of the compositeness parameters.

^{*)} On leave of absence from LPM-USTL Montpellier. Supported in part by the DRCI-CNRS within the France-USA exchange programme: AI-USA 95-03-12.

CERN-TH.4700/87
PM87/19
April 1987

One of the most important experiments planned at LEP I and SLC is the precise measurement of the Z mass in the aim to reach some high accuracy tests of the electroweak theory. Numerous efforts have been investigated for the evaluation of loop effects on the W and Z boson mass shifts within the standard model¹⁾. In this paper, we evaluate various possible leading non-standard contributions induced by the "nearby" compositeness assumption of bosons and fermions on the γ , Z and W propagators. Therefore, we will study successively the effects due to composite (pseudo-) scalars, excited fermions, anomalous spin-one boson couplings and particles with higher spins (2,3/2) to the transverse two-point correlator in the chiral preon symmetry limit:

$$\begin{aligned}\Pi_{ij}^{\mu\nu} &\equiv i \int d^4x e^{iqx} \langle 0 | T J_i^\mu(x) (J_j^\nu(0))^+ | 0 \rangle \\ &= -(g^{\mu\nu} - q^\mu q^\nu / q^2) \Pi_{ij}(q^2),\end{aligned}\quad (1)$$

where $J_i^\mu(x)$ are the operators of bosons which can be written in terms of the particles running in the loop. The normalization of the currents $J_i^\mu(x)$ is done in such a way that the standard vector and axial currents $V^\mu \equiv \bar{\psi}\gamma^\mu\psi$ and $A^\mu \equiv \bar{\psi}\gamma^\mu\gamma^5\psi$ formed with spin- $\frac{1}{2}$ quark bilinears give:

$$J_i^\mu = \lambda_i (a_i V^\mu - b_i A^\mu) \quad (2a)$$

with

$$\lambda_\gamma = g \sin \theta_W \quad ; \quad \lambda_W = g / \sqrt{2} \quad ; \quad \lambda_Z = g / (4 \cos \theta_W)$$

and, for instance, for the up quark:

$$\begin{aligned}a_\gamma &= 2/3 \quad ; \quad a_W = b_W = 1 \quad , \\ a_Z &= 1 - (8/3) \sin^2 \theta_W \quad ; \quad b_Z = 1 \quad .\end{aligned}\quad (2b)$$

Within such a normalization, one can, for instance, write the Z-boson mass shift as¹⁾:

$$\begin{aligned}\frac{\delta M_Z}{M_Z} &= \frac{1}{2} \left\{ \left[\Pi_{WW}(q^2) + \Pi_{\gamma\gamma}(q^2) + \frac{\cos \theta_W}{\sin \theta_W} (1 - t_\theta^2) \Pi_{\gamma Z}(q^2) \right] \frac{1}{q^2} \right\}_{q^2=0} \\ &\quad - \frac{1}{M_Z^2} \text{Re} \Pi_{ZZ}(-M_Z^2) \quad .\end{aligned}\quad (3)$$

We shall evaluate the invariants $\Pi_{ij}(q^2)$ appearing in Eq. (3) by using dimensional regularization²⁾ which is known to preserve gauge invariance (here the photon which is massless is still a gauge field!) and chiral invariance. In fact, we do not have here (in contrast with the old cut-off regularization scheme) any quadratic divergences which spoil the above-mentioned invariances. The most severe divergences which will appear (in a one-loop calculation which we

shall do) are of the type $(M^2)^{-1/\epsilon}$ poles (M^2 is the mass of the particle inside the loop). These divergences will automatically disappear if we have a gauge-type form of the effective Lagrangian (see the following case of scalars). In fact, in some other cases where we have arbitrary couplings, we shall require the disappearance of these leading divergences in order to constrain the arbitrariness for the choice of these couplings.

Equation (3) will serve as phenomenological applications of our calculations. Bearing in mind that a 1% accuracy can be reached at LEP I and SLC for the measurement of the Z mass and assuming that the predictions of the standard model including radiative corrections will agree with the data within this accuracy, we shall therefore explore the constraints on non-standard effects allowed by the 1% window. We shall also notice that the positive properties of the spectral function $\text{Im}\Pi_{ij}(t)$ associated to the invariant $\Pi_{ij}(q^2)$ will allow us to argue that constraints imposed for each individual contribution to the mass shift will be absolute constraints. This is not the case for the ones obtained from (g-2) which is governed by a three-point function having undefined signs.

1. - COMPOSITE (PSEUDO-) SCALARS

We consider the effects of the $\tilde{\pi}$ loop originated from the effective Lagrangian involving (pseudo-) scalar $\tilde{\pi}$ and spin-one weak bosons A^μ :

$$\mathcal{L}_S = -g_\pi \tilde{\pi} \partial_\mu \tilde{\pi} A^\mu + g_\pi^2 \tilde{\pi}^* \tilde{\pi} A_\mu A^\mu. \quad (4)$$

The ordinary two-point correlator induced by the first term reads, for $n = 4-\epsilon$ space-time:

$$\begin{aligned} \tilde{\Pi}_{\tilde{\pi}}^{(2)}(q^2) = & -\left(\frac{1}{3-\epsilon}\right) \cdot \left(-\frac{g_\pi^2}{16\pi^2}\right) \int_0^1 dx \left\{ 8M_\pi^2 + g^2 [1 - 2x(1-x)(6+\epsilon)] \right\} \\ & \cdot \left(\frac{-g^2 x(1-x) + M_\pi^2}{x^2} \right)^{-\epsilon/2} \cdot \Gamma(\epsilon/2), \end{aligned} \quad (5a)$$

while the tadpole contribution is:

$$\tilde{\Pi}_{\tilde{\pi}}^{Tad} = -\left(\frac{1}{3-\epsilon}\right) \left(\frac{g_\pi^2}{16\pi^2}\right) 8M_\pi^2 \left(\frac{M_\pi^2}{x^2}\right)^{-\epsilon/2} \Gamma(\epsilon/2). \quad (5b)$$

It is easy to see that for the sum of the two contributions, the leading M^2 dependence disappears. That renders finite the physical quantities appearing in Eq. (3). This pedagogical example shows the desired constraints which should satisfy other contact terms in the non-trivial examples which we shall discuss

later on. The ν -scale in Eq. (5) is the ϵ -dimension scale which acts like an ultra-violet cut-off up to which the effective theory calculation makes sense. We can therefore identify it with the compositeness scale. The two interesting limits of Eq. (5) for the q^2 values involved in Eq. (3) are:

$$\begin{aligned} \frac{1}{q^2} \tilde{\pi}(q^2) \Big|_{-q^2 \ll M_\pi^2} &\simeq g_\pi^2 \frac{1}{48\pi^2} \left\{ \log \frac{M_\pi^2}{\nu^2} + 2 + \mathcal{O}\left(\frac{q^2}{M_\pi^2}\right) \right\}, \\ \frac{1}{q^2} \tilde{\pi}(q^2) \Big|_{-q^2 = M_\pi^2} &\simeq g_\pi^2 \frac{1}{48\pi^2} \left\{ \log \frac{M_\pi^2}{\nu^2} + 1 \right\}. \end{aligned} \quad (6a)$$

One should notice that the constant terms are regularization-scheme dependent. Using the value $g_\pi = g_W = e/\sin\theta_W$ obtained from p -like universality arguments³⁾, one can deduce for $\nu = 1$ TeV:

$$\frac{\delta M_Z}{M_Z} \simeq \begin{array}{ll} 10^{-3} & : (M_\pi = 0.5 \text{ TeV}) \\ -6 \cdot 10^{-3} & : (M_\pi \simeq M_Z) \end{array}. \quad (6b)$$

Inversely, if we tolerate at most a $\pm 1\%$ deviation between the measured Z -mass at LEP/SLC and the one predicted by the standard model with radiative corrections included, we obtain for $M_\pi = M_Z$:

$$g_\pi \lesssim 0.3. \quad (6c)$$

We have not considered here the case $2M_\pi < M_Z$ as we expect that the best place for seeing this effect is in the Z width. In fact, the $Z \rightarrow \pi\pi$ width, normalized to the phase space factors, is the same as the effect of one generation of fermions³⁾. Now, we study the loops induced by the effective interactions:

$$\mathcal{L} = i g_{pj\ell} \epsilon^{\mu\nu\rho\sigma} p_{1\rho} p_{2\sigma} V_\mu^j V_\nu^\ell P, \quad (7)$$

between two spin-one V_j and one (pseudo-) scalar P . $\epsilon^{\mu\nu\rho\sigma}$ is the anti-symmetric tensor normalized as: $\epsilon^{\mu\nu\rho\sigma} = -\epsilon_{\mu\nu\rho\sigma}$. The evaluation of the loop with internal V_j and P lines gives:

$$\frac{1}{q^2} \tilde{\pi}(q^2) = g_{pj\ell}^2 \frac{1}{16\pi^2} \int_0^1 dx \mathcal{R}^2 (1 - 2 \log \frac{\mathcal{R}^2}{\nu^2}), \quad (8)$$

where $\mathcal{R}^2 = -q^2 x(1-x) + M_P^2(1-x) + M_j^2 x$. One should notice that the gauge dependent part of the V_j propagator which is proportional to $p_{2\rho} p_{2\sigma}' / M_W^2$ gives a null

contribution as expected^{*)}. We deduce, from Eq. (8):

$$\frac{1}{q^2} \Pi(q^2) \Big|_{-q^2 \ll M_P^2} \simeq g_{Pjk}^2 \frac{M_P^2}{16\pi^2} \left\{ -\log \frac{M_P^2}{v^2} + 1 + \frac{1}{2} \left(\frac{M_j}{M_P} \right)^2 + \left(\frac{M_j}{M_P} \right)^4 \left(\log \frac{M_j^2}{M_P^2} - \frac{1}{2} \right) + \mathcal{O}\left(\frac{q^2}{M_P^2}\right) \right\}, \quad (9a)$$

$$\frac{1}{q^2} \Pi(q^2) \Big|_{-q^2 = M_\pi^2, M_j^2 = 0} \simeq g_{Pjk}^2 \frac{M_P^2}{16\pi^2} \left\{ -\frac{5}{6} \log \frac{M_P^2}{v^2} + \frac{11}{8} \right\}. \quad (9b)$$

The value of the $g_{P\gamma\gamma}$ coupling has been estimated from a preon anomaly⁴⁾ and from QHD sum rules³⁾. It is in the range $2\alpha/F_\pi$ and $2\sin\theta_W/M_Z$ where $F_\pi = (\sqrt{2}G_F)^{-\frac{1}{2}}$ is the electroweak breaking scale. The other couplings $g_{\gamma Z P} \dots$ can be obtained from arguments based on vector boson dominance and on global $SU(2)_W$ symmetry. For instance, the smallest coupling which is a technicolour-type coupling contributes to the mass shift as:

$$(\delta M_Z/M_Z) \simeq \mathcal{O}(10^{-5}) : M_P \simeq M_Z, 0.5 \text{ TeV}. \quad (10)$$

Our results for $M_P = M_Z$ are comparable with that obtained in Ref. 5) after a suitable rescaling of the coupling constant. Equation (10) indicates that (pseudo-) scalar bosons having technicolour-type couplings with vector boson pairs have a negligible effect on the Z-mass shift. For a given value of the P-mass, a 1% accuracy on the determination of M_Z , will imply, for $v = 1 \text{ TeV}$:

$$\left(\frac{1}{\alpha} g_{P\gamma\gamma} \right) \lesssim \begin{array}{l} 0.3 \text{ GeV}^{-1} : (M_P = M_Z) \\ 0.07 \text{ GeV}^{-1} : (M_P = 0.5 \text{ TeV}). \end{array} \quad (11)$$

The values of the coupling in Eq. (11) will imply visible signals of (pseudo-) scalar bosons at the $(\bar{p})p$ -TeV colliders from the study of the $(\bar{p})p \rightarrow Z\gamma$ processes⁶⁾. So the combined information from e^+e^- and $(\bar{p})p$ colliders should improve the constraints on the exact value of $g_{P\gamma\gamma}$.

*) In the following, we shall only deal with the effects due to the $g^{\mu\nu}$ (gauge-independent part) of the W,Z propagator. In fact, constraints obtained from the $q^\mu q^\nu$ part of the propagator will be meaningless because of this gauge freedom. However, such contributions are expected to vanish automatically in a consistent effective theory!

2. - THE SPIN- $\frac{1}{2}$ EXCITED FERMION

Let us now consider the magnetic effective coupling:

$$\mathcal{L}_* = g \left(\frac{M_*}{2v} \right) \cdot \bar{\psi}^* \sigma^{\mu\nu} (a + b\gamma_5) \psi F_{\mu\nu}, \quad (12)$$

where ψ^* , ψ and $F_{\mu\nu}$ are fields associated to the excited fermion, standard fermion and bosons. The effects of the ψ^* to the weak boson propagators are given by the loop involving the ψ^* and ψ fields. The evaluation of the two-point correlator implies:

$$\frac{1}{g^2} \frac{\Pi_*(q^2)}{-q^2 \ll M_*^2} \simeq \frac{1}{2} (|a|^2 + |b|^2) N_c \left(\frac{g^2}{8\pi^2} \right) \left(\frac{M_*}{v} \right)^4 \cdot \left(-\frac{1}{3} \right) \left\{ -1 + \frac{g^2}{3M_*^2} \left(\log \frac{M_*^2}{v^2} + 5 \right) \right\}, \quad (13a)$$

$$\frac{1}{g^2} \frac{\Pi_*(q^2)}{-q^2 = M_*^2} \simeq \frac{1}{2} (|a|^2 + |b|^2) \left(\frac{M_*}{v} \right)^4 \left(\frac{g^2}{24\pi^2} \right) N_c \cdot \left(\frac{1}{3} \right) \cdot \left(\log \frac{M_*^2}{v^2} + \frac{52}{3} \right), \quad (13b)$$

where $N_c = 3(1)$ for quarks (leptons). The contribution of the excited fermion to the Z-mass shift for $v = 1$ TeV is:

$$\frac{\delta M_Z}{M_Z} \simeq N_c \left(\frac{g^2}{4\pi} \right) \cdot \begin{cases} 6 \cdot 10^{-3} & (M_* = 0.5 \text{ TeV}) \\ 3 \cdot 10^{-5} & (M_* = M_Z) \end{cases}. \quad (14)$$

The suppression of the contribution in the case $M_* = M_Z$ is mainly due to the v -dependence of the coupling in Eq. (12). In the case where this coupling has no (M_*/v) term, the contribution would become 70 times larger. One should notice that the effects of the excited fermion to the Z-mass shift is small for $g^2 < 4\pi$. The above constraints can be compared with the ones from the (g-2). For $M_* \approx 0(v)$, the (g-2) constraints are stronger⁷⁾ while for $M_* \ll v$, it is shown that the (g-2) results become regularization-scheme dependent^{7b)}.

3. - ANOMALOUS WEAK BOSON COUPLINGS

In the case where the weak bosons are composite, one can expect to have additional three- and four-boson effective couplings compared to the ones of the standard electroweak theory. The forms of such additional terms are known in the literature⁸⁾. Without loss of generality, we restrict ourselves to the CP conserving couplings parametrized à la Ref. 8b):

$$\begin{aligned} \mathcal{L}^{(3)} = (-ie) \left\{ \delta \left[g^{\rho\mu} g^{\nu\lambda} - g^{\rho\nu} g^{\mu\lambda} \right] + \frac{\lambda}{M_W^2} \left[g^{\mu\nu} (P_1 \cdot g P_2^\rho - \right. \right. \\ \left. P_2 \cdot g P_1^\rho - g^{\rho\lambda} (P_1 \cdot P_2 g^{\mu\lambda} - P_2 \cdot g P_1^\mu) + \right. \\ \left. g^{\rho\mu} (P_1 \cdot P_2 g^{\nu\lambda} - P_2 \cdot g P_1^\nu) + P_1^\rho P_2^\nu g^{\mu\lambda} - P_2^\rho g P_1^\mu \right] + \\ \left. \eta_V \epsilon^{\alpha\mu\nu\rho} (P_2 - P_1)_\alpha \right\} V_\rho W_\mu^+ W_\nu^- , \end{aligned} \quad (15)$$

where $\delta+1$, λ and η are magnetic, quadrupole and anapole moments. $q = p_1+p_2$ is the momentum of the V-boson and p_1, p_2 the one of the $W^\pm$. We will also consider the anomalous CP-conserving coupling among three neutral bosons:

$$\mathcal{L}^{(3)} = V_\mu V_{2\alpha} V_{3\beta} (-ie) \left\{ \eta_1^V \epsilon^{\mu\alpha\beta\rho} P_{2\rho} + \frac{1}{M_Z^2} \eta_2^V \epsilon^{\mu\beta\rho\sigma} q_\alpha q_\rho P_{2\sigma} \right\} . \quad (16)$$

and the four-boson couplings:

$$\begin{aligned} \mathcal{L}_V^{(4)} = (-) \left\{ g_1^2 g_{\alpha\beta} g_{\gamma\delta} Z_\gamma Z_\delta W_\alpha W_\beta + g_2^2 (g_{\alpha\beta} g_{\gamma\delta} + g_{\alpha\delta} g_{\beta\gamma}) \cdot \right. \\ \left. W_\alpha W_\beta W_\gamma W_\delta + g_3^2 (g_{\alpha\beta} g_{\gamma\delta} + g_{\alpha\gamma} g_{\beta\delta} + \right. \\ \left. g_{\alpha\delta} g_{\beta\gamma}) Z_\alpha Z_\beta Z_\gamma Z_\delta \right\} . \end{aligned} \quad (17)$$

The couplings g_i^2 in Eq. (17) could be related to each other if one imposes a $SU(2)_W$ global symmetry⁹⁾. However, in order to have a general result we shall keep the g_i as in Eq. (17). We shall see from the following evaluation of the anapole moment contribution that the contact terms in Eq. (17) are necessary for removing the bad behaviour of the two-point correlator at $q^2 = 0$. In fact, these contact terms will act like the scalar tadpole term present in Section 1.

a) Anapole moment

The contribution of the anapole moment is due to the parity-violating term in Eqs. (15) and (16). These contributions do not interfere with the parity-conserving ones in Eq. (15) because in the two-point correlator we can only form one invariant. The evaluations of the W-point correlator involve WV loops ($V \equiv Z, \gamma$). To leading order:

$$\begin{aligned}\pi_W^{(3)}(-q^2 \ll M_W^2) &\simeq \eta_V^2 \left(\frac{e^2}{16\pi^2} \right) \left(-\frac{2}{\epsilon} + \log \frac{M_W^2}{\nu^2} \right) \left[4(M_W^2 + M_V^2) - q^2 \right], \\ \pi_W^{(3)}(-q^2 = M_W^2) &\simeq \eta_V^2 \left(\frac{e^2}{16\pi^2} \right) \left(-\frac{2}{\epsilon} + \log \frac{M_W^2}{\nu^2} \right) [5M_W^2 + 4M_V^2],\end{aligned}\tag{18}$$

where we have neglected the small W-Z mass difference in the internal logarithms. The relevant tadpole diagrams induced by Eq. (17) are:

$$\pi_W^{(4)} = -\frac{2}{12\pi^2} \left(-\frac{2}{\epsilon} + \log \frac{M_W^2}{\nu^2} \right) \left\{ g_1^2 M_Z^2 + \frac{5}{4} g_2^2 M_W^2 \right\}.\tag{19}$$

In the same way, one can evaluate the V-correlator. One obtains to leading order:

$$\begin{aligned}\pi_Z^{(3)}(-q^2 \ll M_W^2) &\simeq \frac{\alpha}{4\pi} \left(-\frac{2}{\epsilon} + \log \frac{M_W^2}{\nu^2} \right) \left\{ (\eta_1^V)^2 + (g^4/M_Z^4) (\eta_2^V)^2 \right\} M_Z^2 + \\ &\quad \eta_V^2 (8M_W^2 - q^2) \eta_V^2 \left\{.\right.\end{aligned}\tag{20}$$

The effect of the tadpole diagram is:

$$\pi_Z^{(4)} = -\frac{2}{12\pi^2} \left(-\frac{2}{\epsilon} + \log \frac{M_W^2}{\nu^2} \right) \left\{ M_W^2 g_1^2 + \frac{3}{2} M_Z^2 g_3^2 \right\}.\tag{21}$$

The finiteness of the anapole moment contributions to the W, Z propagators at $q^2 = 0$ are obtained by cancellation of the $M_{W,Z}^2 (\log M_W^2/\nu^2 - 2/\epsilon)$ terms appearing in Eqs. (18) and (20). If we require such cancellations for each W and Z propagators like the ones in gauge theories, we obtain the sum rules:

$$\begin{aligned}\left(\frac{\alpha}{\pi} \right) \eta_V^2 4(M_W^2 + M_V^2) &\simeq \frac{2}{3} (M_Z^2 g_1^2 + \frac{5}{4} M_W^2 g_2^2), \\ \left(\frac{\alpha}{\pi} \right) (\eta_1^2 M_Z^2 + 8\eta_V^2 M_W^2) &\simeq \frac{2}{3} (M_W^2 g_1^2 + \frac{3}{2} M_Z^2 g_3^2).\end{aligned}\tag{22}$$

The above zeroth-order constraint becomes stronger if we identify the coefficients of the W and Z masses in each equation. Using Eq. (22), we therefore deduce the sum of $\Pi^{(3)}$ and $\Pi^{(4)}$:

$$\begin{aligned} \frac{1}{q^2} \Pi_W(q^2) &\approx -\frac{\alpha}{4\pi} \eta_V^2 \log \frac{M_W^2}{\nu^2}, \\ \frac{1}{q^2} \Pi_Z(q^2) &\approx -\frac{\alpha}{4\pi} \left(\log \frac{M_W^2}{\nu^2} \right) \left[\eta_V^2 - \frac{q^2}{M_Z^2} (\eta_V^V)^2 \right]. \end{aligned} \quad (23)$$

Using Eq. (23) into Eq. (3), we deduce for a 1% accuracy of $\delta M_Z/M_Z$, the approximate constraint for $\nu > 1$ TeV:

$$2 \eta_V^2 + \frac{1}{2} (\eta_V^V)^2 \leq 0.4, \quad (24a)$$

to which corresponds the weakest bound:

$$|\eta_V| \leq 0.4 \quad ; \quad |\eta_V^V| \leq 0.8. \quad (24b)$$

b) Magnetic and quadrupole moments

The two-point correlators associated to the parity conserving terms of Eq. (15) can be due to the diagonal (δ^2, λ^2) , to the interference between λ and δ ($\lambda\delta$) and to the interference of the λ, δ term with the usual standard three-boson couplings (λ, δ) . The evaluations of the various contributions are quite cumbersome. We perform the algebra using a REDUCE programme. We will only keep the leading-log contribution in our analysis. Therefore, the results obtained in four-dimension for the algebra are sufficient. We obtain^{*)}:

$$\begin{aligned} \Pi(q^2) = & -\left(\frac{\alpha}{24\pi}\right) q^2 \left\{ 6\delta^2 + \lambda^2 \left(9 + 2 \frac{q^2}{M_W^2} - \frac{1}{2} \frac{q^4}{M_W^4} \right) \right. \\ & \left. - 6\lambda(5+2\delta) \left(1 - \frac{q^2}{6M_W^2} \right) + 3\delta \right\} \left(-\log \frac{M_W^2}{\nu^2} \right). \end{aligned} \quad (25)$$

^{*)} One should notice that the constraints obtained in Ref. 10c) are regularization scheme and "gauge" dependents. Also, Ref. 10c) does not consider all possible contributing diagrams in his analysis.

From Eqs. (25) and (3), one can deduce the Z-mass shift. We give the constraints on the set (δ, λ) parameters for a 1% deviation of the Z-mass. These are on the ellipse of the Figure. In addition, we show the constraints obtained from the $(g-2)_\mu^{10)}$:

$$-2 \leq \delta \log(v^2/M_W^2) + \lambda \leq 6, \quad (26)$$

which are inside the two bars. Then, we obtain the two sets of constraints:

$$\begin{aligned} & 2 \leq \lambda \leq 4.5 \quad ; \quad -1 \leq \delta \leq 0.3 \\ \text{and} & \\ & -0.2 \leq \lambda \leq 0.3 \quad ; \quad -0.4 \leq \delta \leq 1.2. \end{aligned} \quad (27a)$$

One can also play with the v and precision of the measurement values for giving stronger constraints on λ and δ . These can be obtained for larger v and for increasing precision measurements. We study such effects for $v = 10$ TeV and $\delta M_Z/M_Z \approx 1\%$. There is no significant deviation from the ellipse shown in the Figure. However, the constraints from the $(g-2)_\mu$ become stronger. Then, we obtain:

$$\begin{aligned} & -0.2 \leq \lambda \leq 0.2 \quad ; \quad -0.3 \leq \delta \leq 0.6, \\ & 3.5 \leq \lambda \leq 4.2 \quad ; \quad -0.5 \leq \delta \leq 0.1. \end{aligned} \quad (27b)$$

The constraints in Eq. (27) are useful because they limit the arbitrariness of λ and δ . In fact, the constraints obtained here can be stronger than the ones expected from precise measurements of the polarized $e^+e^- \rightarrow W^+W^-$ processes⁸⁾.

4. - EXCITED Z' AND ISOSCALAR Y BOSONS

The couplings of the excited Z' and isoscalar Y bosons to W^+W^- can be the standard or (and) the anomalous ones in Eq. (15). For the former, one might expect from a phenomenological duality argument that this coupling behaves like $g.(M_Z/M_{Z'})^n$ where $n > 0$, i.e., this effect in the W-propagator is expected to be negligible if $M_{Z'} \gg M_Z$. The effects of the anomalous couplings in Eq. (15) are suppressed compared to that of the Z because of the $\log(M_Z^2/v^2)$ term which controls this effect relative to the $\log(M_Z^2/v^2)$ term for the case of the W and the Z. Therefore, the constraints obtained from the W and the Z boson anomalous terms automatically suppress the effects of excited and isoscalar bosons. Couplings of the Z' and Y to the neutral gauge bosons are expected to be anomalous. Therefore, they are of the type discussed in Eq. (16). Our discussions on the Z'W⁺W⁻ anomalous couplings also apply here.

5. - PARTICLES WITH HIGHER SPINS

In addition to the above anomalous effects of compositeness, the existence of particles with higher spins should be characteristic consequences of any bound state models. Here we shall limit ourselves to the effects of spin-2 and spin-3/2 particles.

a) Spin-2 bosons

Some possible phenomenological effects of such bosons have already been studied in Refs. 6) and 11). The coupling of this boson to a spin-1 boson pair can be described by the effective vertex:

$$V^{\rho\sigma} = -\frac{i}{v} \left\{ g_2 F_1^{\rho\mu} F_{2\mu}^{\sigma} + \frac{g_1}{\sqrt{2}} F_1^{\mu\nu} F_{2\mu\nu} P_1^{\rho} P_2^{\sigma} \right\}, \quad (28a)$$

and we work in the gauge where the spin-2 propagator is:

$$D^{\rho\sigma\alpha\beta} = \frac{1}{2} \left\{ g^{\rho\alpha} g^{\sigma\beta} + g^{\rho\beta} g^{\sigma\alpha} - \frac{2}{3} g^{\rho\sigma} g^{\alpha\beta} \right\} \frac{1}{q^2 - M_2^2}, \quad (28b)$$

with

$$q = p_1 + p_2.$$

The leading order contributions to the W, Z correlators are given by the coupling involving the lowest dimension operators. In fact, we have checked that the effects of the high dimension coupling are suppressed by $(p_2/v)^2$, where $-p_2^2 < M_W^2$ is the external momentum. Then, we obtain the two-point correlator:

$$\frac{1}{q^2} \Pi(q^2) \simeq g_2^2 \left(\frac{3}{8\pi^2} \right) \left(\frac{M_2}{v} \right)^2 \log \frac{M_2^2}{v^2} + \mathcal{O}(q^2/v^2). \quad (29)$$

M_2 is the mass of the spin-2 boson. A priori, we do not have any idea of the value of this mass. Using a QHD sum rule analysis of the spin-2 correlator, we obtain the local duality constraint¹²⁾:

$$M_2^2 \simeq \frac{3}{4} t_c \left\{ 1 + \frac{80\pi}{27} \alpha_H \langle F^2 \rangle / t_c^2 \right\}, \quad (30)$$

for a bound state of spin- $\frac{1}{2}$ preon and for $SU(3)_H$. This sum rule should be compared to the one of the W-mass¹³⁾:

$$M_W^2 \simeq \frac{t_c}{2} \left\{ 1 - \frac{2\pi}{3} \alpha_H \langle F^2 \rangle / t_c^2 \right\}, \quad (31)$$

where a small value of the W-mass is dual to a large value of the hypergluon condensate $\alpha_H \langle F^2 \rangle \approx 3/2 \pi t_c^2$ if the preon continuum threshold $\sqrt{t_c}$ is of the order of 1 TeV. With these values of the parameters, Eq. (30) implies that the spin-2 mass should be in the TeV range. Therefore, in order to see the numerical effect of Eq. (29) on the Z-mass shift, let us use $(M_2/v) \approx 0.5$. Then, for $\delta M_Z/M_Z \approx 1\%$, we obtain:

$$g_2 \lesssim 0.3 \quad . \quad (32)$$

b) Spin-3/2 excited fermions

Spin-3/2 fermions can be described by the following lowest dimension interpolating fields¹⁴⁾:

$$\psi_\mu^* = (g_{\mu\nu} - \frac{1}{4} \gamma_\mu \gamma_\nu) D^\nu \psi \quad . \quad (33)$$

(D^ν is the covariant derivative and ψ the spin- $\frac{1}{2}$ field) from which one can deduce the spin-3/2 propagator. The effective fermion-fermion gauge interaction is:

$$\mathcal{L} = \frac{1}{v} \left\{ g_{1/2} \bar{\psi}_\mu^* \psi A^\mu + \epsilon^{\mu\nu\rho\sigma} \bar{\psi}_\mu^* \gamma_5 \gamma_\nu \psi A_\rho \right\} \quad . \quad (34)$$

The contribution of the first type of coupling to the boson propagator is:

$$\tilde{\Pi}_{1/2} = g_{1/2}^2 \left(\frac{M_*}{v} \right)^2 \left(\log \frac{M^2}{v^2} \right) \frac{1}{4\pi^2} \left(-\frac{5}{3} q^2 + 4 M_*^2 \right), \quad (35)$$

where we have used the REDUCE programme for the algebra. The second term in Eq. (35) is still disappointing. Its presence reflects the incomplete form of the exotic Lagrangian. The effects of the second term in Eq. (34) are high dimension in the mass dependence. To leading order:

$$\tilde{\Pi}_{3/2}(q^2) = g_{3/2}^2 \frac{11}{8\pi^2} M_*^6 \log \frac{M_*^2}{v^2} \quad . \quad (36)$$

One should note that due to the positive properties of the two-point spectral function, there is no way of eliminating the leading M^2 term by adding Eqs. (35) and (36). Assuming that such unwanted terms disappear in a much more complete theory, the lowest-leading dimension effects on the weak boson masses are due to the q^2 -term in Eq. (35). This effect is very similar in strength with the one in Eq. (29) and leads to:

$$|g_{1/2}| \lesssim 0.3 \quad . \quad (37)$$

6. - FERMION FORM FACTORS AND CONCLUSIONS

We have postponed here the discussion on the effects due to the fermion form factors. Such effects can be deduced from the point-like fermion ones by a rescaling factor: $(1 \pm q^2/v^2)$ where q is the momentum of the external weak boson. Using the fact that the contributions of point-like fermions are of the order of $3 \text{ GeV}^{15)}$, one might expect an effect of form factors less than 25 MeV for v larger than 1 TeV. So, the effects are negligible. This discussion closes our systematic analysis of the nearby compositeness effects on the electroweak boson propagators. We have applied the results of our calculation on the Z-boson mass-shift value. The extension of these applications to some other electroweak processes can certainly be done. We plan to come back to these points in forthcoming work.

ACKNOWLEDGEMENTS

We thank H. Fritzsch, B. Lynn, A. Pich, E. de Rafael, D. Schildknecht, G. Veneziano and especially M. Perrottet for discussions. We acknowledge the kind hospitality of the CERN Theoretical Physics Division.

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FIGURE

Behaviour of the sets (λ, δ) for given values of $\delta M_Z/M_Z$ and ν . The ellipse is the line where (λ, δ) should stay, as given by the result deduced from Eq. (25). The hatched lines are the frontiers of the domain permitted by $(g-2)_\mu$ for different values of the compositeness scale.