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MASSIVE YANG MILLS - AN EFFECTIVE LAGRANGIAN FOR COMPOSITE $W^{\pm}$, Z AND NEW COLORED VECTOR BOSONS

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Abstract: The framework is weak interactions, interpreted as residual (hypercolor) interactions among composite q,l,$W^{\pm}$ and Z. An effective Lagrangian $\mathcal{L}_{eff}$ for "low energies" ($E \lesssim G_F^{-1/2}$) is derived from i) a specification of the global chiral symmetry G of weak interactions (for $\alpha, \alpha_{color} \to 0$), ii) local $U(1)_{em} \times SU(3)_c$ gauge invariance and iii) vector boson dominance in the operator form of current-field identities. The result is a massive Yang-Mills Lagrangian with respect to the global group G.

$\mathcal{L}_{eff}$ for q,l,W,Z interactions, basing on $G = SU(2)_{WI}$ of global weak isospin, turns out to be identical (in its dimension ≤ 4 operator part) to the Lagrangian of the standard (GSW) model in the unitary gauge without the physical Higgs. $\mathcal{L}_{eff}$ predictions are argued to closely mimic the GSW predictions due to the chiral nature of G and the smallness of the effective coupling constant.

An extension of the scheme to larger symmetry groups as expected from preon models for $\alpha, \alpha_{color} \to 0$ (e.g. $G = SU(2)_{WI} \times SU(4)_{Pati-Salam}$) is proposed. This implies the existence of new colored (and uncolored) composite vector bosons and vector dominance in the gluon sector. $\mathcal{L}_{eff}$ then determines the interactions of these new bosons with quarks and leptons in terms of a few free parameters. Interesting consequences for $p\bar{p}$ collider and HERA experiments as well as for precision experiments at low energies emerge.

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1. INTRODUCTION

The Glashow-Salam-Weinberg (GSW) model [1] of electroweak interactions describes charged and neutral current reactions with ease and predicted the W and Z vector bosons with correct masses. Nevertheless, the confirmation of its "hardcore", the renormalizable local gauge theory nature, is lacking up to now: the Higgs scalar has not (yet) been found and its indirect influence through tree and loop contributions is so far undetectably small for a Higgs mass ranging from a few GeV up to order 1 TeV.

This fact leaves room for the conceptually quite different idea [2)-6)] that quarks and leptons as well as the W and Z vector bosons could be composite. This interpretation, originally based on purely theoretical motivation, has received further stimulation from some recent experimental hints [7),8)] for possible deviations from the GSW predictions at the CERN $p\bar{p}$ collider.

In a composite picture of this type one usually assumes, in analogy to QCD, an underlying confining hypercolor gauge theory on the preon level with compositeness scale

$$\Lambda \sim G_F^{-1/2} \sim 300 \text{ GeV}. \tag{1}$$

Weak interactions in the presently explored energy range, $E \lesssim m_{W,Z}$, then qualitatively parallel strong interactions for $E \lesssim m_\rho$: they appear as short-range, "low-energy" residual hypercolor interactions among the hypercolor singlet, composite quarks, leptons, W and Z vector bosons. However, unlike strong interactions, a chiral protection mechanism à la 't Hooft [9)] is required, in order to keep the quark and lepton masses small as compared to Λ and to reconcile [10)] a scale as low as $G_F^{-1/2}$ with $(g-2)_\mu$ measurements [11)].

Let us emphasize that in such a composite scenario the W and Z bosons are not related to gauge bosons of a renormalizable local gauge theory. They just represent prominent composite vector bosons, analogous to the ρ mesons in strong interactions. As a characteristic signature of compositeness one would expect a (possibly rich) spectrum of further composite states, bosons and fermions, somewhere between 100 Gev and 1 TeV.

The theoretical interest ultimately focusses of course on the underlying renormalizable, local hypercolor gauge theory on the preon level. This is reflected in the intensive activity in composite model building.

However, experiments are so far restricted to the deep infrared region ($E \lesssim m_{W,Z}$) of the (hypothetical) hypercolor theory. Thus, it is of considerable phenomenological interest to acquire some understanding of the effective "low-energy" interactions among hypercolor singlet composites like q, l, $W^{\pm}$ and Z, e.g. in form of an effective Lagrangian description. Such an investigation serves a two-fold purpose. First of all, one hopes to understand why the low-energy interactions of composite q, l, $W^{\pm}$ and Z should mimic so closely the well-established GSW interactions. Secondly, one looks for predictions for further composites (presumably heavier than the W and Z bosons) and their effective interactions with quarks and leptons. Such predictions are badly needed in order to distinguish the "composite picture" from the "elementary GSW picture". Differences could show up already for $E \lesssim m_{W,Z}$ in precision measurements or of course directly and more dramatically for $E \gtrsim m_{W,Z}$ by the appearance of new states. These issues are addressed in this paper.

A promising route towards a description of the effective "low-energy" interactions among "old" (q, l, $W^{\pm}$ and Z) and possible new composites is, to transfer to the regime of weak interactions well-known concepts which have successfully described low-energy strong interactions among composite hadrons (well before the advent of QCD).

Following such a strategy, principles like vector-boson-photon mixing (Bjorken [12]; Hung and Sakurai [13]), vector boson dominance (Kögerler and Schildknecht [14]) and current algebra (Fritzsch, Kögerler and Schildknecht [15]) have been transferred to weak interactions with considerable success. W-boson dominance has led [14] to explain the observed structure of the neutral current in the composite framework as well as to correctly predict (!) $m_{W,Z}$ in terms of G_F and $\sin\Theta_W$. SU(2)-current algebra naturally explains [15] the universality of the couplings of composite W,Z bosons to quark and lepton pairs.

The aim of this paper is a systematic effective Lagrangian approach for weak interactions in the "low-energy" regime, $E \lesssim \Lambda \sim G_F^{-1/2} \sim 300$ GeV . This effective Lagrangian will imply the full wisdom which may be abstracted and generalized from low energy strong interactions: current algebra in the strongest realization as field algebra, Weinberg's sum rules and in particular vector boson dominance in the operator formulation as current-field identities

$$\text{conserved currents} \propto \text{composite vector boson fields} \qquad (2)$$

An important new aspect of this approach is a natural generalization of vector boson dominance to the gluon sector, predicting among others also new colored composite vector bosons and their effective interactions with quarks and leptons.

The main characteristic of the resulting effective Lagrangian is its massive Yang-Mills structure with respect to the global symmetry group of weak interactions.

The paper is organized as follows. In Sect. 2 we formulate the three requirements from which the effective Lagrangian ($\mathcal{L}_{eff}$) is to be constructed: i) a specification of the global chiral symmetry G for α, $\alpha_{color} \to 0$, ii) local $U(1)_{em} \times SU(3)_c$ gauge invariance and iii) the powerful requirement of current-field identities for the corresponding local symmetry currents, involving a multiplet of composite vector bosons in the adjoint representation of G. We start in Sect. 3 by recapitulating the classical construction of $\mathcal{L}_{eff}$ from the analogous input ingredients for the prototype case of strong interactions, as presented by Lee and Zumino [16] in 1967. Sect. 4 is devoted to the simplest application of our program to weak interactions involving only the known particles, q, l, $W^{\pm}$, Z, and basing on $G = SU(2)_{WI}$ of global weak isospin (for $\alpha \to 0$). The result is an effective Lagrangian of the massive Yang-Mills type which is formally identical in its dimension ≤ 4 operator part to the GSW Lagrangian in the unitary gauge without physical Higgs. The implications of this result are discussed in detail. In particular, we argue that $\mathcal{L}_{eff}$ predictions closely mimic GSW predictions due to the chiral nature of the global symmetry and the smallness of the universal effective coupling constant. In Sects. 5 and 6 we then extend the scheme to larger global symmetries of the weak interactions, entailing the presence of new uncolored and colored composite vector bosons. Sect. 6, in particular, is devoted to the interesting case of a current-field identity for the color octet currents involving an octet of new colored vector bosons. As an illustration we explicitely construct $\mathcal{L}_{eff}$ corresponding to an enlarged global symmetry (for α and $\alpha_{color} \to 0$) as expected in a popular class of preon models : $G = SU(2)_{WI} \times SU(4)_{Pati-Salam}$. The interactions of the new vector bosons with quarks and leptons are determined in terms of two mass and two coupling parameters. Interesting consequences for $\bar{p}p$-collider and HERA experiments as well as for precision measurements of G_F, $\sin\Theta_W$, m_W and m_Z emerge. Sect.7 contains a summary and conclusions.

2. SYMMETRIES AND CURRENT FIELD IDENTITY

In this Section we propose the systematic exploitation of two concepts in the framework of effective interactions of composite quarks, leptons, W,Z bosons etc.:

i) the symmetry content as emerging for $\alpha, \alpha_c = 0$ and $\alpha, \alpha_c \neq 0$ and

ii) the powerful principle of current-field identities as abstracted from strong interactions.

(i) Symmetry content: At distances of order $G_F^{1/2}$ the color gauge coupling is small

$$\alpha_c(G_F^{-1/2}) = \frac{g_c^2(G_F^{-1/2})}{4\pi} \sim 0.1 \tag{3}$$

due to asymptotic freedom, which allows to consider the color gauge interactions along with the electromagnetic gauge interactions as soft perturbations of the effective weak interactions. In the limit $\alpha, \alpha_c \to 0$, weak interactions will have a certain unbroken global symmetry. The global symmetry group G will have to contain

$$G \supset SU(2)_{WI}, \text{ the global } SU(2) \text{ of weak isospin} \tag{4}$$

and

$$G \supset U(1)_{em} \times SU(3)_c; \tag{5}$$

furthermore, for consistency G must be chiral, in order to implement 't Hooft's chiral protection mechanism [9] on the composite level, i.e. to keep the composite quarks and leptons massless (on the scale $\Lambda \sim G_F^{-1/2}$). This property is an important difference to strong interactions. It will be at the root of the surprising success of our prescription in weak interactions (as compared to strong interactions).

The (unknown) global symmetry G is the most important link to the underlying preon theory. In fact, on the low-energy composite level, the only manifestation of the preon theory is through its characteristic global symmetry G and the classification of the ground state composite spectrum with respect to G. Thus preon models, with 't Hooft's anomaly constraints incorporated, will serve as a guide for the choice of a specific (unbroken) global group G.

Switching on the gauge couplings α and α_c breaks G explicitely and softly down to the local gauge symmetry $U(1)_{em} \times SU(3)_c$.

(ii) Current-field identities: First, a massive composite vector boson is associated with each global symmetry current, i.e. the global symmetry G is assumed to prescribe the spectrum of prominent vector bosons such that the vector boson multiplet transforms as the adjoint representation of G.

The powerful dynamical requirement comes in for $\alpha, \alpha_c \neq 0$: a current-field identity for each local symmetry current, in our case the electromagnetic current and the color octet currents, is required to hold. For the electromagnetic current this is the operator formulation of vector-boson dominance as familiar from strong interactions, for the color octet currents it is a generalization to a "vector-boson dominance for gluons".

Following and generalizing the logics of the beautiful paper [16] by Lee and Zumino, written in 1967 in the context of strong interactions, we shall cast these two principles in the form of an effective, "low-energy" Lagrangian for weak interactions. This $\mathcal{L}_{eff}$ will imply the full wisdom [16),17] which may be abstracted and generalized from low-energy strong interactions: current algebra in the strongest realization as field algebra with respect to the global group G, Weinberg's sum rules and vector-boson dominance with respect to the photons as well as the gluons. For illustration see Figs. 1a,b.

This analysis opens the door to relate the global chiral symmetry as abstracted from preon models to the spectrum of "old" ($W^{\pm}$, Z) and "new" (e.g. colored) composite vector bosons and their low-energy interactions with composite quarks and leptons. It predicts small deviations from GSW even at energies $\lesssim m_Z$ in terms of a few parameters (masses and couplings) which may allow to soon pin down the global symmetry content of weak interactions (in the limit $\alpha, \alpha_c \to 0$). This in turn would strongly constrain composite model building on the preon level.

3. A REMINDER: CURRENT-FIELD IDENTITY AND EFFECTIVE LAGRANGIAN IN STRONG INTERACTIONS

As a reminder of the power of the principle of current field identity let us first return to the familiar framework of strong interactions and briefly recapitulate the Lee-Zumino derivation [16] of the effective Lagrangian for the simplest example. Input is the global SU(2) isospin symmetry which is exact for $\alpha \to 0$, a triplet of composite ρ-vector meson fields $\vec{\rho}_\mu$, with mass m (for $\alpha \to 0$), local $U(1)_{em}$ gauge invariance for $\alpha \neq 0$ and the current-field identity for the I = 1 component of the electromagnetic current

$$j_\mu^{em}\Big|_{I=1} = \frac{m^2}{g}\,\rho_\mu^3 \qquad \text{with } Q = T_3 + Y \tag{6}$$

This means the conserved electromagnetic current is chosen as an interpolating field for the (composite) uncharged ρ meson. At this input level, the proportionality constant m^2/g, or equivalently g, is simply an unknown constant.

An effective Lagrangian for $\alpha = 0$ a priori admits a large variety of couplings in terms of the ρ-meson fields. The combination of three conditions [16] on $\vec{\rho}_\mu(x)$, however, strongly restricts these couplings. The first one comes from the equation of motion for $\vec{\rho}_\mu(x)$, the second one from the conservation of the global SU(2) symmetry currents. The third one provides the most powerful restriction: the immediate consequence of the conservation of the electromagnetic current, the current-field identity (6) and global SU(2) symmetry in the limit $\alpha \to 0$ is the field conservation equation

$$\partial^\mu \vec{\rho}_\mu(x) = 0 \tag{7}$$

(corresponding to the spin 1 condition for a massive vector field). The striking result [16], for $m \neq 0$, is an effective Lagrangian for the strong interactions ($\alpha = 0$) of the massive Yang-Mills type

$$\mathcal{L}_{eff}^{\alpha=0} = \mathcal{L}_{Y\text{-}M} + \frac{1}{2} m^2 \vec{\rho}_\mu \vec{\rho}^\mu \tag{8}$$

which can be generalized to include any further hadron fields to

$$\mathcal{L}_{eff}^{\alpha=0} = \mathcal{L}_{Y\text{-}M}(\vec{\rho}_{\mu\nu}, \text{fields}, \mathcal{D}_\mu \text{ fields}) + \frac{1}{2} m^2 \vec{\rho}_\mu \vec{\rho}^\mu \tag{9}$$

$\mathcal{L}_{Y\text{-}M}$ is the most general (non-renormalizable) Lagrangian exhibiting local SU(2) isospin gauge invariance, where the triplet of ρ vector meson fields plays the rôle of the gauge fields and the constant g, defined in eq. (6), the rôle of the universal gauge coupling constant. One has to keep in mind, however, that the ρ mesons are really composite hadrons and the coupling g is really an effective coupling. The tensor $\vec{\rho}_{\mu\nu}$ and $\mathcal{D}_\mu$ are the familiar non-abelian field-strength

tensor and covariant derivative, respectively. The only term breaking the local SU(2) gauge invariance down to a global one is the ρ-meson mass term! The only way how any further hadron fields, like e.g. the nucleon isodoublet field, can couple to the ρ-meson field is through $\mathcal{D}_\mu$, involving the single coupling constant g (hence ρ universality), and through $\vec{\rho}_{\mu\nu}$.

When switching on $\alpha \neq 0$, the combined effect of the requirements of local $U(1)_{em}$ gauge invariance and of the current-field identity (6) results [16] in the following prescription. $\mathcal{L}_{eff}^{\alpha\neq 0}$ is obtained from $\mathcal{L}_{eff}^{\alpha=0}$ by replacing $\vec{\rho}_\mu$ in $\mathcal{L}_{Y\text{-}M}$, but not in the ρ-mass term, by $\hat{\vec{\rho}}_\mu$ defined as

$$\hat{\rho}_\mu^{1,2} = \rho_\mu^{1,2}, \qquad \hat{\rho}_\mu^3 = \rho_\mu^3 + \frac{e}{g} A_\mu \tag{10}$$

and by adding the appropriate kinetic term of the photon. Obviously, the massive Yang-Mills structure is retained in the case $\alpha \neq 0$ in terms of the fields $\hat{\vec{\rho}}_\mu$. For further details and generalizations we refer to Ref. 16.

To summarize, even though initially only global SU(2) symmetry of strong interactions ($\alpha = 0$) was required, local $U(1)_{em}$ gauge invariance ($Q = T_3 + Y$) and the current-field identity (6) for the symmetry current corresponding to the T_3 generator, enforce the massive Yang-Mills structure of the strong interaction Lagrangian. Of course $\mathcal{L}_{eff}$ is non-renormalizable and its region of applicability is limited to sufficiently low energies ($E \lesssim O(m_\rho)$).

In strong interactions the effective Lagrangian satisfying the current-field identity has essentially only been of esthetical value. First of all, the relevant effective coupling constant is very large, $g = g_\rho \sim 5.5$, i.e.

$$g_\rho^2/4\pi \sim 2.4 > 1. \tag{11}$$

Furthermore, related to this, the approximation by single ρ-meson exchange barely makes sense, since it violates the unitarity bound at energies closely above the ρ-meson mass. As is well known, in strong interactions the unitarity bound is taken care of by excitations of new hadronic states close to m_ρ ("reggeization"). Correspondingly it is not surprising that vector meson dominance and an evaluation of current algebra in terms of a single ρ meson are typically violated on the 10-20% level in strong interactions.

In weak interactions we shall find a much more favorable situation for the practical applicability of $\mathcal{L}_{eff}$ due to two properties

i) the effective coupling constant is known [14),15)] to be small:

$$g^2 / 4\pi \sim 0.03 \ll 1 , \tag{12}$$

ii) the chiral symmetry is (approximately) unbroken, keeping the leptons and quarks approximately massless.

This will be discussed at the end of Sect. 4.

4. AN EFFECTIVE LAGRANGIAN FOR COMPOSITE QUARKS, LEPTONS AND $W^{\pm}$, Z BOSONS

As pointed out already in the introduction, two important steps towards an understanding of the effective "low-energy" interactions among composite q, l, $W^{\pm}$ and Z had been taken in Refs. 14,15. In Ref. 14 the idea of W-dominance was explored in the neutral current sector, whereas in Ref. 15 universality of the $Wf\bar{f}$ couplings emerged as a consequence of $SU(2)_{WI}$-current algebra and W-pole dominance. The effective Lagrangian obtained below will automatically imply the results of both, Ref. 14 and Ref. 15, make them more transparent and will lead beyond them.

In this Section we specialize to the simplest application of the general program outlined in Sect. 2. We just concentrate on the $SU(2)_{WI}$-weak isospin part of the global symmetry G and a corresponding triplet of composite vector boson fields $\vec{w}_\mu$ with mass m_W (for $\alpha \to 0$). These fields are to represent the well-known vector bosons of weak interactions. For the time being, we disregard any (presumably heavier) vector bosons associated with the $U(1)_Y$ or $SU(2)_R \times U(1)_{B-L}$ subgroups or with the enlarged global symmetry for α and $\alpha_c \to 0$. This idealization is consistently possible, if for α, $\alpha_c \to 0$ the global symmetry group G is non-simple, i.e. of the form

$$G = SU(2)_{WI} \times G' . \tag{13}$$

The results of this Section then refer to the limit, where the masses of the composite vector bosons associated with the G' symmetry tend to infinity. We shall return to the more realistic and notably more interesting case of finite masses for these bosons in the following two Sections.

The composite fermion fields f (generically for quarks and leptons) are known to transform with respect to the global $SU(2)_{WI}$ as follows: the left-handed ones (f_L) as doublets, the right-handed ones (f_R) as singlets. This chiral fermion classification properly implements parity violation in weak interactions and at the same time suppresses quark and lepton mass terms, just as required in the composite framework! The fermions f, moreover, have to realize a global $U(1)_Y$ symmetry with standard hypercharge assignments $y_{L,R}$. The corresponding conserved current is denoted by j_μ^Y.

Now, with the (composite) fields $\vec{w}_\mu$ and f (=q,l) specified, we are ready to construct the effective Lagrangian from the requirements of

α) global $SU(2)_{WI} \times U(1)_Y$ invariance for $\alpha = 0$

β) local $U(1)_{em}$ gauge invariance for $\alpha \neq 0$ and

γ) the current-field identity for the weak isovector part of the electromagnetic current

$$j_\mu^{em} = j_\mu^{T_3} + j_\mu^Y \tag{14}$$

$$j_\mu^{T_3} = \frac{m_W^2}{g_W} w_\mu^3 \tag{15}$$

The construction of $\mathcal{L}_{eff}$ runs parallel to our discussion of the strong interaction example in Sect. 3, basing on Ref. 16. Again, for $\alpha = 0$, $\mathcal{L}_{eff}^{\alpha=0}$ has the form of a massive $SU(2)_{WI}$-Yang-Mills Lagrangian, with $\vec{w}_\mu$ playing the rôle of the "quasi-gauge" fields and g_W the rôle of the universal "quasi-gauge" coupling. Again, $\mathcal{L}_{eff}^{\alpha\neq 0}$ for $\alpha \neq 0$ is obtained i) by replacing everywhere in $\mathcal{L}_{eff}^{\alpha=0}$, except in the vector boson mass term, $\vec{w}_\mu$ by $\vec{W}_\mu$, with

$$W_\mu^{1,2} = w_\mu^{1,2} , \qquad W_\mu^3 = w_\mu^3 + \frac{e}{g_W} A_\mu , \tag{16}$$

ii) by introducing a standard gauge invariant coupling of the photon to the hypercharge in the covariant derivative such that altogether

$$\mathcal{D}_\mu^{\alpha=0} = \partial_\mu + i g_W \vec{w}_\mu \vec{T} \tag{17}$$

$$\longrightarrow \mathcal{D}_\mu^{\alpha\neq 0} = \partial_\mu + i g_W \vec{W}_\mu \vec{T} + ie A_\mu Y$$

$$(= \partial_\mu + i g_W \vec{w}_\mu \vec{T} + ie A_\mu Q)$$

and finally iii) by adding the appropriately normalized kinetic term for the photon field A_μ. In Eq. (17), $\vec{T}$, Y and $Q = T_3 + Y$ are the generators of $SU(2)_{WI}$, $U(1)_Y$ and $U(1)_{em}$, respectively. They have to be taken in the matrix representation appropriate for the field $\mathcal{D}_\mu$ is acting on.

The resulting effective Lagrangian in terms of the input fields $\vec{w}_\mu$, f and A_μ reads (cf. Eqs.(8,9))

$$\mathcal{L}_{eff}^{\alpha\neq 0} = \mathcal{L}_{Y-M}^{\alpha\neq 0} + \frac{1}{2} m_W^2 \vec{w}_\mu \vec{w}^\mu \tag{18}$$

with

$$\begin{aligned} \mathcal{L}_{Y-M}^{\alpha\neq 0} = & -\frac{1}{4} F_{\mu\nu} F^{\mu\nu} \left(1-\left(\frac{e}{g_W}\right)^2\right) - \frac{1}{4} \vec{W}_{\mu\nu} \vec{W}^{\mu\nu} \\ & + i \bar{f}_L \gamma_\mu \mathcal{D}^\mu f_L + i \bar{f}_R \gamma_\mu \mathcal{D}^\mu f_R \\ & + \mathcal{L}'_{dim>4} (\vec{W}_{\mu\nu}; f_L, \mathcal{D}_\mu f_L; f_R, \mathcal{D}_\mu f_R) \end{aligned} \tag{19}$$

Here

$$\begin{aligned} F_{\mu\nu} &= \partial_\mu A_\nu - \partial_\nu A_\mu \\ \vec{W}_{\mu\nu} &= \partial_\mu \vec{W}_\nu - \partial_\nu \vec{W}_\mu - g_W \vec{W}_\mu \times \vec{W}_\nu \\ \mathcal{D}_\mu f_L &= (\partial_\mu + i g_W \vec{W}_\mu \vec{\tau}/2 + i e A_\mu y_L) f_L \\ \mathcal{D}_\mu f_R &= (\partial_\mu + i e A_\mu y_R) f_R . \end{aligned} \tag{20}$$

$\vec{W}_\mu$ is given in terms of $\vec{w}_\mu$ and A_μ by Eq.(16). Summation over the fermions f is implied. The factor $(1 - (e/g_W)^2)$ attached to $-1/4\, F_{\mu\nu} F^{\mu\nu}$ guarantees the proper normalization of the photon kinetic term, since there is a further contribution, $-1/4\, (e/g_W)^2 F_{\mu\nu} F^{\mu\nu}$ coming from the $-1/4\, \vec{W}_{\mu\nu} \vec{W}^{\mu\nu}$ term.

$\mathcal{L}'_{dim>4}$ in Eq. (19) contains the non-renormalizable contributions to $\mathcal{L}_{Y-M}$, i.e. all possible operators of dimension > 4 composed of $\vec{W}_{\mu\nu}$, f_L, $\mathcal{D}_\mu f_L$, f_R and $\mathcal{D}_\mu f_R$ in a Lorentz-invariant, locally $SU(2)_{WI}$ gauge invariant and locally $U(1)_{em}$ gauge invariant way.

Let us next discuss the precise symmetry content of the various contributions to $\mathcal{L}_{eff}^{\alpha\neq 0}$, then demonstrate that $\mathcal{L}_{eff}^{\alpha\neq 0}$ indeed implies the current-field identity (14,15) and finally confront $\mathcal{L}_{eff}^{\alpha\neq 0}$ with the GSW Lagrangian.

By inspection one realizes that the essential part of the effective Lagrangian, namely $\mathcal{L}_{Y-M}^{\alpha\neq 0}$, Eq. (19), exhibits a <u>local</u> $SU(2)_{WI} \times U(1)_Y$ gauge invariance, if $\vec{W}_\mu$ (instead of $\vec{w}_\mu$) and

$$B_\mu = A_\mu \sqrt{1-(e/g_W)^2} \tag{21}$$

are considered as independent gauge fields, i.e.

$$\delta \mathcal{L}_{Y-M}^{\alpha\neq 0} = 0 \tag{22}$$

for infinitesimal local gauge transformations

$$\delta f = i\,(\vec{\alpha}(x)\vec{T} + \beta(x) Y) f \tag{23}$$

$$\delta \vec{W}_\mu = -\frac{1}{g_W} \partial_\mu \vec{\alpha}(x) - \vec{\alpha}(x) \times \vec{W}_\mu \tag{24}$$

$$\delta B_\mu = \delta A_\mu \sqrt{1-(e/g_W)^2} = -\frac{\sqrt{1-(e/g_W)^2}}{e} \partial_\mu \beta(x) . \tag{25}$$

Let us again emphasize that the <u>local</u> gauge nature is <u>not</u> an input, but the consequence of conditions α)- γ). $U(1)_{em}$ is readily identified as the subgroup of the local $SU(2)_{WI} \times U(1)_Y$ generated by

$$\vec{\alpha}(x) = (0, 0, \beta(x)) \tag{26}$$

in Eqs. (23,24), corresponding to $Q = T_3 + Y$.

The local $SU(2)_{WI} \times U(1)_Y$ symmetry is only broken by the $\vec{w}_\mu$ mass term, such that the full effective Lagrangian $\mathcal{L}_{eff}^{\alpha\neq 0}$ exhibits the required local $U(1)_{em}$ gauge invariance. This is easiliy verified: the $U(1)_{em}$ transformations (23)-(26) induce the following change of the fields $\vec{w}_\mu$ appearing in the mass term

$$\delta \vec{w}_\mu = \delta \vec{W}_\mu - (0, 0, \frac{e}{g_W} \delta A_\mu) = -\vec{\alpha}(x) \times \vec{w}_\mu , \tag{27}$$

with $\vec{\alpha}(x)$ as in Eq. (26). This transformation law, appropriate for a massive $SU(2)_{WI}$ triplet, of course leaves the mass term invariant, since the $\partial_\mu \vec{\alpha}(x)$ terms have cancelled.

The current-field identity is obtained from $\mathcal{L}_{eff}^{\alpha\neq 0}$ by means of the field equations for the original fields A_μ and w_μ^3

$$\frac{\delta \mathcal{L}_{eff}^{\alpha\neq0}}{\delta A_\mu} = \frac{\delta \mathcal{L}_{Y-M}^{\alpha\neq0}}{\delta A_\mu} = 0 \qquad \text{for } w_\mu^3 \text{ fixed} \tag{28}$$

and

$$\frac{\delta \mathcal{L}_{eff}^{\alpha\neq0}}{\delta w_\mu^3} = \frac{\delta \mathcal{L}_{Y-M}^{\alpha\neq0}}{\delta w_\mu^3} + m_W^2 w^{3,\mu} = 0 \quad \text{for } A_\mu \text{ fixed} \tag{29}$$

where the variational derivative is defined as usual

$$\frac{\delta \mathcal{L}}{\delta\, \text{field}} = \frac{\partial \mathcal{L}}{\partial\, \text{field}} - \partial_\nu \frac{\partial \mathcal{L}}{\partial(\partial_\nu \text{field})} \tag{30}$$

Next, we remember that apart from the kinetic term, A_μ enters in $\mathcal{L}_{Y-M}^{\alpha\neq0}$ only through

$$W_\mu^3 = w_\mu^3 + \frac{e}{g_W} A_\mu \tag{31}$$

and

$$(\partial_\mu + ie A_\mu y_{L\atop R}) f_{L\atop R} \tag{32}$$

in the covariant derivative. Thus

$$\frac{\delta \mathcal{L}_{Y-M}^{\alpha\neq0}}{\delta w_\mu^3} = \frac{\delta \mathcal{L}_{Y-M}^{\alpha\neq0}}{\delta W_\mu^3} \qquad \text{for } A_\mu \text{ fixed} \tag{33}$$

and

$$\frac{\delta \mathcal{L}_{Y-M}^{\alpha\neq0}}{\delta A_\mu} = \frac{\delta \mathcal{L}_{Y-M}^{\alpha\neq0}}{\delta W_\mu^3}\frac{e}{g_W} + \partial_\nu F^{\nu\mu}\left(1-\left(\frac{e}{g_W}\right)^2\right) - e\, j^{Y,\mu} \tag{34}$$

for w_μ^3 fixed

with

$$j_\mu^Y = -i\left[\frac{\partial \mathcal{L}_{eff}^{\alpha\neq0}}{\partial(\partial^\mu f_L)} y_L f_L + \frac{\partial \mathcal{L}_{eff}^{\alpha\neq0}}{\partial(\partial^\mu f_R)} y_R f_R\right] = \bar f_L \gamma_\mu y_L f_L + \bar f_R \gamma_\mu y_R f_R + \dots \tag{35}$$

being the conserved hypercharge current.

Upon combining Eqs. (28,29,33,34), we obtain the following form⁺ of the

⁺ The formalism of Lee and Zumino [16] refers to leading order in e/g, as appropriate for strong interactions where $e/g \ll 1$. Our effective Lagrangian (18-20) and current-field identity (36) represent the straightforward generalization to include all orders in e/g.

current-field identity (14,15)

$$\partial^\nu F_{\nu\mu} = \tilde e\, j_\mu^{em} = \frac{e}{1-(e/g_W)^2}\left(\frac{m_W^2}{g_W} w_\mu^3 + j_\mu^Y\right) \tag{36}$$

For consistency we need

$$\frac{e}{g_W} < 1 \tag{37}$$

which is conform with our input notion that electromagnetic gauge interactions break the global $SU(2)_{WI}$ of weak interactions softly.

After the discussion of the symmetry properties of $\mathcal{L}_{eff}^{\alpha\neq0}$ the main result of this Section does not come as a surprise: $\mathcal{L}_{eff}^{\alpha\neq0}$ is formally identical on the level of the dimension ≤ 4 operators to the GSW Lagrangian in the unitary gauge without the physical Higgs field, provided we identify

$$\frac{e}{g_W} \equiv \sin\Theta_W, \quad \text{i.e.} \quad g_W = g_{GSW} \sim 0.64 \tag{38}$$

and

$$A_\mu = B_\mu / \cos\Theta_W \tag{39}$$

(see also Eq. (21)). In terms of the fields $\vec W_\mu$ and B_μ the effective Lagrangian $\mathcal{L}_{eff}^{\alpha\neq0}$, Eqs. (18-20), takes the familiar GSW form (without Higgs)

$$\begin{aligned}\mathcal{L}_{eff}^{\alpha\neq0} = &-\tfrac14 B_{\mu\nu}B^{\mu\nu} - \tfrac14 \vec W_{\mu\nu}\vec W^{\mu\nu} \\ &+ \tfrac12 m_W^2 \{(W_\mu^1)^2 + (W_\mu^2)^2\} + \tfrac12 m_W^2 (W_\mu^3 - \mathrm{tg}\Theta_W B_\mu)^2 \\ &+ i\bar f_L \gamma^\mu(\partial_\mu + i g_W \vec W_\mu \tfrac{\vec\tau}{2} + i\frac{e}{\cos\Theta_W} B_\mu y_L) f_L \\ &+ i\bar f_R \gamma^\mu(\partial_\mu + i\frac{e}{\cos\Theta_W} B_\mu y_R) f_R + \mathcal{L}'_{dim>4}\end{aligned} \tag{40}$$

The "current-mixing term"

$$-\frac14 \sin\Theta_W (F_{\mu\nu} w^{3,\mu\nu} + w_{\mu\nu}^3 F^{\mu\nu}) \tag{41}$$

in the basis w_μ^3, A_μ or equivalently the "mass-mixing term"

$$-\frac12 m_W^2 \mathrm{tg}\Theta_W (W_\mu^3 B^\mu + B_\mu W^{3,\mu}) \tag{42}$$

in the basis W_μ^3, B_μ can be removed by the (familiar) linear field transformation to the diagonalized basis

$$A_\mu = A_\mu^{diag} - tg\,\theta_W Z_\mu \qquad\qquad B_\mu = \cos\theta_W A_\mu^{diag} - \sin\theta_W Z_\mu$$
$$W_\mu^3 = Z_\mu/\cos\theta_W \quad (43a) \qquad W_\mu^3 = \sin\theta_W A_\mu^{diag} + \cos\theta_W Z_\mu \quad (43b)$$

which implies the finite mass renormalization due to electromagnetic mixing

$$m_Z = m_W/\cos\theta_W \qquad (44)$$

From Eq. (43a) we see that the massless field A_μ^{diag} is related to the massless input field A_μ by a canonical transformation. It is, however, clear that the important dynamical issue, the current-field identity (36), is transparent only if expressed in terms of A_μ and not of A_μ^{diag}.

For completeness let us quote the current-field identity (36) in terms of the fields Z_μ and B_μ

$$\partial^\nu B_{\nu\mu} = e\left(\frac{m_Z^2}{g_W} Z_\mu + \frac{1}{\cos\theta_W} j_\mu^Y\right). \qquad (45)$$

What can one conclude from our result that the effective massive Yang-Mills Lagrangian on the leading (dimension ≤ 4) operator level is identical to the GSW Lagrangian in the unitary gauge without physical Higgs?

In comparison with a massive Yang-Mills Theory the hardcore of the GSW theory lies in the presence of the Higgs scalar which turns it into a renormalizable local gauge theory. However, the Higgs particle has not (yet) been found and (so far) Higgs contributions to experimentally accessible observables are undetectably small[18] on the tree and one-loop level and also on the two-loop level as far as calculated.

Of course one could imagine that besides the composite q, l, $W^\pm$ and Z there is a composite scalar with all the properties of the Higgs scalar, thus turning $\mathcal{L}_{eff}^{\alpha\neq0}$ into a renormalizable Lagrangian of the GSW type. This situation has been envisaged and discussed[+] in Ref. 19.

[+] We learnt about this work in preparation, when presenting the results of this paper at the "Workshop on Quark and Lepton Structure" in Erice/Sicily, April 1984

On the other hand, there is just one known example of effective interactions among composites resulting from an underlying, confining gauge theory: the strong interactions among hadrons. They suggest the following alternative which we tentatively have adopted. The effective Lagrangian (of the massive Yang-Mills type) for composite q, l, $W^\pm$ and Z is indeed non-renormalizable and correspondingly has a limited range of validity. The tree diagrams calculated from it will eventually come into conflict with the unitarity bound (which goes hand in hand with non-renormalizability [20]). This conflict is resolved by reggeization, i.e. by the appearance of a rich spectrum of excited composites.

On the basis of such an analogy to strong interactions, the Higgs scalar may be considered as a clean signal for the GSW theory and the appearance of new (composite) states as a signal for compositeness.

As long as neither the Higgs nor new composites show us the way, it is important to ask to which extent the massive Yang-Mills Lagrangian can mimic the GSW predictions which agree so well with the data.

First of all (to the extent that contributions from $\mathcal{L}'$ in Eq. (19) may be neglected) we expect

$$\text{tree level results from } \mathcal{L}_{eff}^{\alpha\neq0} \approx \text{tree level results from } \mathcal{L}_{GSW}. \qquad (46)$$

This general statement implies known results [+ 14),15)] concerning the tree-level predictions for the neutral current, for m_W and [++] for $m_Z = m_W / \cos\Theta_W$ and for the universality of the $Wf\bar{f}$ couplings. It of course implies many more, e.g.

i) triple and quartic W,Z couplings $\approx$ corresponding GSW couplings

ii) the gyromagnetic ratio g of the $W^\pm$: $g(W^\pm) \approx 2$ (= GSW value)[+++]

[+] In Ref. 14), precisely speaking, W-dominance combined with the abelian global symmetry $U(1)_{T_3} \times U(1)_Y$ was exploited.

[++] The corresponding relation for strong interactions, $m_{\omega} = m_{\rho^\pm} / \cos\Theta_\rho$ with $\sin\Theta_\rho = e/g_\rho$, $g_\rho \sim 5.5$ is compatible with experiment: $m_{\rho^\pm} = 766.7 \pm 2.8$ MeV, $m_{\rho^\pm} / \cos\Theta_\rho = 768.1 \pm 2.8$ MeV to be compared with $m_{\rho^0} = 769.7 \pm .86$ MeV. Unfortunately the errors are too large to make this a meaningful statement.

[+++] The gyromagnetic ratio of the ρ meson is still not measured. In the context of composite W and Z it would be interesting to see whether the Lee-Zumino prediction [16], $g(\rho^\pm) \approx 2$, holds true also for the ρ meson.

iii) tree level partial width for processes like

$W,Z \longrightarrow q\bar{q},\ l\bar{l},\ q\bar{q}q\bar{q},\ l\bar{l}l\bar{l},\ q\bar{q}l\bar{l},\ \ldots \approx$

corresponding GSW partial widths,

etc.

The quality of the approximation (46) increases with increasing distance between $m_{W,Z}$ and the threshold m^* for the excited composite spectrum (W', spin two boson,... This threshold will in turn roughly coincide with the energy where the tree diagrams get into conflict with unitarity bounds.

In this respect, the situation is much more favorable than in strong interactions. We recall [20),21)] that in a massive Yang-Mills theory with unbroken chiral symmetry, i.e. with $m_f/m_W \longrightarrow 0$, all $2 \longrightarrow 2$ tree amplitudes, except those involving three or four longitudinally polarized vector bosons, satisfy the tree unitarity constraints on the power level for $E \longrightarrow \infty$ (ignoring log E effects). All tree unitarity violations are pushed to high energies due to the smallness of the effective coupling g_W and to powers of $m_f/m_W \ll 1$ in amplitudes involving fermions. Thus the smallness of g_W combined with chiral symmetry leads one to expect a large gap between $m_{W,Z}$ and m^*, typically m^* between a few 100 GeV and 1 TeV (see also Ref. 22 in this context). Further support comes from duality arguments [14),23)]. Thus, altogether, for energies $E < m^*$ relation (46) should be a good approximation.

This has to be contrasted with strong interactions, where the effective coupling constant is large ($g_\rho \sim O(10)\ g_W$) and chiral symmetry is spontaneously broken ($m_N > m_\rho$), leading in fact to $m^* \sim m_{\rho'} \sim m_{f\text{-meson}} \sim 1.5\ m_\rho$, i.e. m^* very close to m_ρ.

In this context, let us point out that Visnjić has recently suggested [24)] a causal relation between the size of the effective coupling g and the realization of chiral symmetry, by associating the chiral low-energy weak four fermion interaction with a Nambu-Jona-Lasinio type interaction. As a result [24)] small couplings $g < g_{crit}$ are argued to be necessary for unbroken chiral symmetry, as appropriate for the composite weak interaction scenario, whereas for $g > g_{crit}$ spontaneous breakdown occurs, consistent with the situation in strong interactions.

The massive Yang-Mills Lagrangian being non-renormalizable, one cannot make any predictions on the one-loop level. However, it is instructive to recall [25)] that, again in the chiral limit, $m_f/m_W \longrightarrow 0$, the dimension ≤ 4 operator part of the massive Yang-Mills Lagrangian is in fact almost[+] one-loop renormalizable.

[+] in the sense that one-loop renormalizability is only violated by a logarithmic divergence associated with the four W vertex.

Of course one has no quantitative control over contributions from $\mathcal{L}'$ (in Eq. (19)) and/or from excited composites (even if one were to introduce a cut-off of the order of m^*, say, which is justifiable in the composite framework). However, it would not come as a surprise, if the one-loop contributions, calculated for $m_f \longrightarrow 0$ from the dimension ≤ 4 operator part of $\mathcal{L}_{eff}$, were to account for the major contribution beyond the tree level.

A suggestive example is the weak interaction contribution to the anomalous magnetic moment of the μ : $\Delta(g-2)_\mu$. This quantity is tightly bounded [11)] from above by experiment and QED calculations. It tends to represent a hurdle for the composite interpretation of weak interactions. Naive dimensional counting [10)], in the presence of chiral symmetry, leads to

$$\Delta (g-2)_\mu \sim O(1) \left(\frac{m_\mu}{\Lambda}\right)^2 \tag{47}$$

where $O(1)$ reflects one's ignorance about the "one-loop coefficient". The experimental bound implies [11)] $\Lambda \gtrsim 700$ GeV which is only marginally consistent with the compositeness interpretation pursued in this paper.

On the other hand, the GSW one-loop prediction [26)] is

$$\Delta_{GSW} (g-2)_\mu = \frac{g_W^2}{4\pi^2}\left(\frac{2}{3}\sin^4\Theta_W - \frac{1}{3}\sin^2\Theta_W + \frac{1}{4}\right)\left(\frac{m_\mu}{m_W}\right)^2 \times \left(1 + O\left(\left(\frac{m_\mu}{m_\phi}\right)^2\right)\right). \tag{48}$$

It is finite (after renormalization) in the unitary gauge without the Higgs scalar, the Higgs contribution itself being negligibly small, of order $(m_\mu/m_W)^4$. In fact the original calculation by Jackiw and Weinberg [26)] was performed in the unitary gauge. Largely due to the small size of g_W, $g_W^2/4\pi \sim .03$, this contribution to $\Delta(g-2)_\mu$ is an order of magnitude below the present experimental sensitivity. The one-loop contribution to $\Delta(g-2)_\mu$ from the dimension ≤ 4 operator part of $\mathcal{L}_{eff}$ in the limit $m_f \longrightarrow 0$ will be approximately equal to the GSW contribution; thus there is still room of an order of magnitude for further contributions over which we have no quantitative control.

Possibly, the message from this exercise is that in a weak-coupling effective interaction, where $g^2/4\pi$ is not of order 1 but very small as compared to 1, one should replace the estimate (47) by

$$\Delta(g-2)_\mu \sim O(1)\ \frac{g_W^2}{4\pi}\left(\frac{m_\mu}{\Lambda}\right)^2 \tag{49}$$

leading to $\Lambda \gtrsim$ 120 GeV. This in turn is perfectly consistent with a compositeness interpretation of the weak interactions.

In conclusion, the effective "low-energy" interactions among composite q, l, $W^\pm$ and Z, resulting from our massive Yang-Mills Lagrangian, closely mimic the GSW interactions. The excited composite spectrum is only expected well above $m_{W,Z}$.

Of course one expects new ground state composites with masses possibly even close to $m_{W,Z}$. They all have to be appropriately included in $\mathcal{L}_{eff}$. This applies in particular to the new composite (uncolored and colored) vector bosons to be discussed in the next two Sections. It also applies, e.g. to the composite isoscalar spin 0 boson [27] ($m \sim$ 50 GeV) or the composite massive lepton [28] ($m \sim$ 80 GeV) which have been proposed in the context of the radiative $Z \to e^+e^-\gamma$ events observed at the CERN $\bar{p}p$ collider [7),8)].

Their couplings are severely constraint by the following requirements

i) chiral symmetry, keeping m_f small,

ii) the massive Yang-Mills structure of $\mathcal{L}_{eff}$,

iii) their contribution to $\Delta(g-2)_\mu$ has to respect the experimental bound.

5. NEW UNCOLORED COMPOSITE VECTOR BOSONS

The effective Lagrangian involving only the known particles, q, l, $W^\pm$ and Z as composites has been shown to successfully mimic GSW predictions and thus also the data. This gives us confidence for the next step, the inclusion of new (massive) composite vector bosons, associated with the global symmetry of weak interactions larger than $SU(2)_{WI}$. This will lead to interesting deviations from the GSW predictions even at energies lower than the new boson masses.

In this Section we address the issue of one or more additional vector bosons related to the global $U(1)_Y$ symmetry of weak interactions (cf. also Ref. 14). There is an ambiguity as to which is the correct full global symmetry for $\alpha \to 0$. Possibilities are $SU(2)_{WI} \times \tilde{G}$ with

global group $\tilde{G}$	Y generator	associated composite vector boson fields
$U(1)_Y$	Y	y_μ
$U(1)^L_Y \times U(1)^R_Y$	$Y = Y_L + Y_R$	$y^L_\mu,\ y^R_\mu$
$SU(2)_R \times U(1)_{B-L}$	$Y = T^3_R + \frac{B-L}{2}$	$\vec{w}^R_\mu,\ u_\mu$
$SU(2)_R \times U(1)^L_{B-L} \times U(1)^R_{B-L}$	$Y = T^3_R + (\frac{B-L}{2})_L + (\frac{B-L}{2})_R$	$\vec{w}^R_\mu,\ u^L_\mu,\ u^R_\mu$

(50)

where the chiral U(1) symmetries seem in fact more appropriate for the case of (approximately) massless quarks and leptons.

Clearly each simple factor in the global group introduces a vector boson multiplet in the adjoint representation and two new free parameters, a mass m and a "quasi-gauge" coupling g. Depending on the composition of the charge generator $Q = T^3_{WI} + Y$ in terms of the available generators (see Eq.(50)) the current-field identity will contain contributions from a different set of massive vector bosons

$$\begin{aligned} j^{em}_\mu &= \frac{m_W^2}{g_W} W^3_\mu + \frac{m_Y^2}{g_Y} y_\mu \\ &\vdots \\ j^{em}_\mu &= \frac{m_W^2}{g_W} W^3_\mu + \frac{m^{R\,2}_W}{g^R_W} W^{R,3}_\mu + \frac{m^{L\,2}_u}{g^L_u} u^L_\mu + \frac{m^{R\,2}_u}{g^R_u} u^R_\mu \end{aligned} \tag{51}$$

The recipe for constructing the effective Lagrangians implementing the respective current-field identities and local $U(1)_{em}$ gauge invariance is the same as applied for $SU(2)_{WI}$. The simplest case of implementing the $U(1)_Y$ has been treated in Ref 14. The effective Lagrangian in terms of the fields $\vec{w}_\mu$, y_μ, A_μ and f, exhibiting local $U(1)_{em}$ gauge invariance, the full global symmetry $SU(2)_{WI} \times U(1)_Y$ for $\alpha \to 0$ and satisfying the current-field identity (51) has the following form

$$\mathcal{L}_{eff}^{\alpha\neq 0} = -\frac{1}{4} F_{\mu\nu}F^{\mu\nu}\left(1-\left(\frac{e}{g_W}\right)^2-\left(\frac{e}{g_Y}\right)^2\right) - \frac{1}{4}\vec{W}_{\mu\nu}\vec{W}^{\mu\nu} - \frac{1}{4}Y_{\mu\nu}Y^{\mu\nu}$$
$$+\frac{1}{2}m_W^2\,\vec{w}_\mu\vec{w}^\mu + \frac{1}{2}m_Y^2\,y_\mu y^\mu \tag{52}$$
$$+ i\bar{f}_L\gamma^\mu\mathcal{D}_\mu f_L + i\bar{f}_R\gamma^\mu\mathcal{D}_\mu f_R + \mathcal{L}'_{dim>4}(\vec{W}_{\mu\nu}, Y_{\mu\nu}, f_L, \mathcal{D}_\mu^L f_L, f_R, \mathcal{D}_\mu^R f_R)$$

$\vec{W}_\mu$ and $\vec{W}_{\mu\nu}$ are defined as in Eqs. (16) and (20), respectively; analogously we have

$$Y_\mu = y_\mu + \frac{e}{g_Y}A_\mu$$
$$Y_{\mu\nu} = \partial_\mu Y_\nu - \partial_\nu Y_\mu \tag{53}$$

Furthermore

$$\mathcal{D}_\mu = \partial_\mu + i g_W \vec{W}_\mu\vec{T} + i g_Y Y_\mu Y \tag{54}$$

where the generators $\vec{T}$ of $SU(2)_{WI}$ and Y of $U(1)_Y$ are to be taken in the respective representations of the field $f_{L,R}$ they are acting on. Again, all the "current-mixing" terms $-1/4\, e/g_W\,(w^3_{\mu\nu}F^{\mu\nu} + F_{\mu\nu}w^{3,\mu\nu})$ and $-1/4\, e/g_Y\,(y_{\mu\nu}F^{\mu\nu} + F_{\mu\nu}y^{\mu\nu})$ can be removed by an appropriate linear field transformation from the basis A_μ, w^3_μ and y_μ to the basis A_μ^{diag}, Z_μ and y_μ^{diag}. This basis change, again a canonical one with respect to the photon field, implies a finite mass renormalization, worked out in detail in Ref. 14. We quote it here for later reference

$$\left.\begin{matrix} m^2(Z) \\ m^2(y^{diag}) \end{matrix}\right\} = \frac{1}{2}\,\frac{1}{1-s_W^2-s_Y^2}\left[c_Y^2 m_W^2 + c_W^2 m_Y^2 \mp\sqrt{(c_W^2 m_Y^2 - c_Y^2 m_W^2)^2 + (2 s_W s_Y m_W m_Y)^2}\right] \tag{55}$$

$$\text{with}\quad s_W = \frac{e}{g_W},\; s_Y = \frac{e}{g_Y};\; c_W = \sqrt{1-s_W^2},\; c_Y = \sqrt{1-s_Y^2}.$$

The same mass-mixing is obtained for $U(1)_Y^L \times U(1)_Y^R$, if y_μ is associated with $U(1)_Y^L$ and $m_Y^R \to \infty$. (The only changes occur in the current-field identity and in the definition of $\mathcal{D}_\mu$).

6. CURRENT-FIELD IDENTITY FOR GLUONS AND NEW COLORED COMPOSITE VECTOR BOSONS

In Sects. 4,5 we restricted ourselves to the framework defined by the global symmetry of weak interactions in the limit $\alpha \to 0$, local $U(1)_{em}$ gauge invariance for $\alpha \neq 0$ and current-field identity for the conserved electromagnetic current. Next we extend the framework to the enlarged (unknown) unbroken global symmetry G of weak interactions in the limit $\alpha, \alpha_c \to 0$, local $U(1)_{em} \times SU(3)_c$ gauge invariance for $\alpha, \alpha_c \neq 0$ and current-field identities for the conserved electromagnetic and color octet currents.

Following the strategy outlined in Sect.2 we propose the existence of a composite vector boson multiplet transforming like the adjoint representation of G. Since necessarily $G \supset SU(3)_c$, there will be a color octet of composite vector bosons, v_μ^a, $a = 1,\ldots,8$, among them. Next we use the (covariantly) conserved color octet currents as interpolating fields for these color octet vector bosons, i.e. we require the current-field identities (in analogy to Eq. (36))

$$(\mathcal{D}^\nu_{color})^{ab}\, G^b_{\nu\mu} = \frac{g_c}{1-\left(\frac{g_c}{g_V}\right)^2}\,\frac{m_V^2}{g_V}\,v_\mu^a + \ldots \qquad a = 1,\ldots 8, \tag{56}$$

where $(\mathcal{D}^\nu_{color})^{ab}\, G^b_{\nu\mu} \propto$ the covariantly conserved color octet currents (summation over b being implied). $\mathcal{D}^\nu_{color}$ is the covariant $SU(3)_c$ derivative, $G^a_{\nu\mu}$ the gluon field strength tensor and $g_c = g_c(\Lambda)$ the running coupling constant at the hypercolor scale $\Lambda \sim G_F^{-1/2} \sim 300$ GeV. This is the straightforward generalization of vector-boson dominance to the case of gluons +, as depicted in Fig. 1b. m_V is the mass of the color octet vector bosons in the limit $\alpha, \alpha_c \to 0$, g_V their effective weak interaction coupling constant.

The whole picture is only consistent for

$$g_c/g_V < 1 \qquad (g_c = g_c(G_F^{-1/2})) \tag{57}$$

which is analogous to condition (37). This inequality implies that the color

+ When presenting the results of this paper at the "Workshop on Quark and Lepton Structure" in Erice/Sicily, April 1984, we learnt about related work in preparation by W. Buchmüller[29] and also by U. Baur and K.H. Streng[30] and by D. Düsedau, D. Lüst and D. Zeppenfeld[31]. W. Buchmüller has independently explored colored vector bosons associated with a global symmetry similar to the one considered in this paper.

gauge interactions have to be weaker than the effective residual weak interactions. It specifies the notion that color gauge interactions are considered to break the global symmetry G <u>softly</u>.

The next step is the construction à la Lee and Zumino [16] of the effective "low-energy" Lagrangian. It involves a generalization from the <u>abelian</u> local gauge symmetry ($U(1)_{em}$) to a non-abelian one ($SU(3)_c$) which turns out to go through without complication.

Let us present the construction for a representative example. The generalization to any other global group G is then quite obvious. The example is abstracted from a whole class of what we phrase "Abbott-Farhi type" preon models, Refs. 4,6,32.

Let us briefly recapitulate the essentials of this class of preon models. The underlying confining hypercolor gauge symmetry on the preon level is $SU(2)_{HC}$. The preon content consists of a doublet scalar ϕ, giving rise to the global $SU(2)_{WI}$ (for $\alpha \to 0$) and four left-handed, massless fermions F (for one family), giving rise to a global, <u>chiral</u> $SU(4)^L$ symmetry (for $\alpha, \alpha_c \longrightarrow 0$).

$$SU(4)^L \supset SU(3)_c^L \times U(1)_{B-L}^L \tag{58}$$

is a left-handed Pati-Salam SU(4) and of course

$$\left(\frac{B-L}{2}\right)_L = Y_L \;, \quad Q_L = T^3_{WI} + Y_L \,. \tag{59}$$

In the left-handed sector the global symmetry is altogether

$$G_L = SU(2)_{WI} \times SU(4)^L \tag{60}$$

with

$$F = (\underline{1}, \underline{4}) \quad \text{where} \quad \underline{4} = \underset{\substack{\uparrow \\ Q}}{\underline{3}}{}_{\frac{1}{6}} + \underset{\substack{\uparrow \\ L}}{\underline{1}}{}_{-\frac{1}{2}} \quad \text{with respect to (58)}$$
$$\phi = (\underline{2}, \underline{1}) \tag{61}$$

where Q denotes the color triplet and L the color singlet preons in the multiplet F.

The left-handed quarks and leptons are bound states of the type

$$f_L = \phi^\dagger F \,, \qquad \text{i.e.} \quad q_L = \phi^\dagger Q \,, \; \ell_L = \phi^\dagger L \,. \tag{62}$$

They transform like $(\underline{2}, \underline{4})$ with respect to G_L, Eq. (60), and are kept massless, consistent with 't Hooft's anomaly conditions.

In this class of models the familiar $SU(2)_{WI}$ triplet of composite W bosons is associated with the $SU(2)_{WI}$ currents on the preon level

$$SU(2)_{WI}: \qquad W_\mu^i \longleftrightarrow i\phi^\dagger \frac{\tau^i}{2} \mathcal{D}_\mu^{HC} \phi + h.c. \,, \quad i = 1, 2, 3, \tag{63}$$

where $\mathcal{D}_\mu^{HC}$ is the covariant derivative of the underlying hypercolor $SU(2)_{HC}$ and τ^i the $SU(2)_{WI}$ matrices. Correspondingly, we associate with the global $SU(4)^L$ currents a 15-plet of <u>new</u> composite vector bosons

$$SU(4)^L: \qquad v_\mu^A \longleftrightarrow \bar{F} \frac{\lambda^A}{2} \gamma_\mu F \,, \qquad A = 1,\ldots,15 \tag{64}$$

where $\lambda^A/2$ are the $SU(4)^L$ matrices. Summation over hypercolor indices is implied in Eqs. (63,64). The $SU(3)_c^L \times U(1)_{B-L}^L$ decomposition of v_μ^A is as follows

$$\underline{15} = \underset{\substack{\uparrow \\ v_8}}{\underline{8}}{}_0 + \underset{\substack{\uparrow \\ v_3}}{\underline{3}}{}_{4/3} + \underset{\substack{\uparrow \\ \bar{v}_3}}{\underline{\bar{3}}}{}_{-4/3} + \underset{\substack{\uparrow \\ v_1 = y}}{\underline{1}}{}_0 \tag{65}$$

where we label the vector bosons by their color content and realize that v_1 is identical to y (more precisely to y^L), introduced in Sect.5. We then have

$$\begin{aligned} v_8 &\longleftrightarrow \bar{Q} \frac{\lambda^a}{2} \gamma_\mu Q \,, \quad a = 1, \ldots 8 \\ v_3 &\longleftrightarrow \frac{1}{\sqrt{2}} \bar{L} \gamma_\mu Q \\ y &\longleftrightarrow \sqrt{\frac{3}{2}} \left(\frac{1}{6} \bar{Q} \gamma_\mu Q - \frac{1}{2} \bar{L} \gamma_\mu L \right), \end{aligned} \tag{66}$$

where $\lambda^a/2$ for a = 1,...,8 are identified with the $SU(3)_c^L$ matrices.

As concerns the right-handed quarks and leptons, they are either pointlike [4),32)] or composite [6)] with residual hypercolor interactions mediated by vector bosons much heavier than those of the left-handed sector. In any case, to first approximation justice is done to all variants, if all composite vector bosons associated with the right-handed sector are considered to be infinitely massive. The right-handed fermions then only experience $SU(3)_c \times U(1)_{em}$ gauge interactions. All what remains of relevance for our purposes is that the global symmetry G_R of the right-handed fermions

$$G_R \supset SU(3)^R_c \times U(1)^R_{em} \qquad (\text{with } Q^R = Y^R) \tag{67}$$

with

$$\begin{aligned} SU(3)_c &= \text{diagonal subgroup of } SU(3)^L_c \times SU(3)^R_c \\ U(1)_{em} &= \text{diagonal subgroup of } U(1)^L_{em} \times U(1)^R_{em} \end{aligned} \tag{68}$$

such that color and electromagnetic gauge interactions become vectorlike.

Next, we forget about the details of the underlying preon models and construct the effective Lagrangian from the following requirements

i) <u>global</u> $G = G_L \times G_R$-symmetry for $\alpha, \alpha_c \longrightarrow 0$, where

$$G_L = SU(2)_{WI} \times SU(4)^L \tag{69}$$

$$\left.\begin{aligned} &\supset SU(3)^L_c \times U(1)^L_{em} \\ G_R &\supset SU(3)^R_c \times U(1)^R_{em} \end{aligned}\right\} \supset \left[SU(3)_c \times U(1)_{em}\right]_{L+R} \tag{70}$$

and G is to be realized by the following set of composite fields

$$\begin{aligned} &\text{left-handed fermions} && f_L = (q,1)_L = \begin{cases} (\underline{2},\underline{4}) & \text{of } G_L, \\ \text{singlets of } G_R \end{cases} \\ &\text{right-handed fermions} && f_R = (q,1)_R\,, \quad \text{singlets of } G_L \end{aligned} \tag{71}$$

and transforming as usual under $SU(3)_c \times U(1)_{em}$, and

composite vector bosons in the <u>adjoint</u> of G_L, singlets of G_R, i.e.

$$\vec{w}_\mu = (\underline{3},\underline{1}) \quad \text{of mass } m_W \quad (\text{for } \alpha \to 0), \tag{72}$$

$$\begin{aligned} \vec{v}_\mu &= (\underline{1},\underline{15}) \quad \text{of mass } m_V \quad (\text{for } \alpha, \alpha_c \longrightarrow 0) \\ &= v_8 + v_3 + \bar{v}_3 + y \end{aligned} \tag{73}$$

$\vec{w}_\mu$ are associated with the $SU(2)_{WI}$ currents as in Sect.4, the 15-plet $\vec{v}_\mu$ of <u>new</u> vector bosons with the $SU(4)^L$ currents;

ii) local $SU(3)_c \times U(1)_{em}$ gauge invariance for $\alpha, \alpha_c \neq 0$ with A_μ and G^a_μ denoting the photon and gluon gauge fields, respectively;

iii) the current-field identities (cf. Eqs. (36,51,56))

$$\partial^\nu F_{\nu\mu} = \frac{e}{1-\left(\frac{e}{g_W}\right)^2 - \frac{2}{3}\left(\frac{e}{g_V}\right)^2} \left[\frac{m_W^2}{g_W} w^3_\mu + \frac{m_V^2}{g_V}\sqrt{\frac{2}{3}}\, v^{15}_\mu + j^{Y,R}_\mu\right] \tag{74}$$

$$\left(\partial^\nu_{color} G_{\nu\mu}\right)^a = \frac{g_c}{1-\left(\frac{g_c}{g_V}\right)^2} \left[\frac{m_V^2}{g_V} v^a_\mu + j^{a,R}_{\mu,color}\right], \quad a = 1,\ldots,8. \tag{75}$$

In Eq. (74), v_μ^{15} corresponds to the y vector boson associated with the hypercharge $U(1)^L_Y \subset SU(4)^L$ ($= y_\mu^L$ in Eq. (50)). SU(4) symmetry gives for the hypercharge matrix

$$y_L = \sqrt{\frac{2}{3}}\, \frac{\lambda^{15}}{2} \tag{76}$$

and accordingly

$$g_Y = \sqrt{\frac{3}{2}}\, g_V. \tag{77}$$

Moreover, in Eq. (75), v_μ^a denote the components of the composite color octet vector boson v_8 contained in the 15-plet $\vec{v}_\mu$ (cf. Eq. (65)).

The effective Lagrangian satisfying i)-iii) is of the massive Yang-Mills type with respect to the global input symmetry G_L, Eq. (69). It contains altogether only <u>four free parameters</u>, a vector boson mass and a quasi-gauge coupling for each simple factor in G_L: m_W, g_W for $SU(2)_{WI}$ and m_V, g_V for $SU(4)^L$.

By straightforward generalization of Eqs. (18,19) and Eq. (52) one finds the following result

$$\begin{aligned} \mathcal{L}^{\alpha,\alpha_c \neq 0}_{eff} = \mathcal{L}_{Y-M} &+ \frac{1}{2} m_W^2\, w^i_\mu\, w^{i,\mu} \\ &+ \frac{1}{2} m_V^2\, v^A_\mu\, v^{A,\mu} \end{aligned} \tag{78}$$

with

$$\mathcal{L}_{Y-M} = -\frac{1}{4} F_{\mu\nu}F^{\mu\nu}\left(1-\left(\frac{e}{g_W}\right)^2-\frac{2}{3}\left(\frac{e}{g_V}\right)^2\right) - \frac{1}{4} G^a_{\mu\nu}G^{a,\mu\nu}\left(1-\left(\frac{g_c}{g_V}\right)^2\right)$$
$$-\frac{1}{4} W^i_{\mu\nu}W^{i,\mu\nu} \quad - \frac{1}{4} V^A_{\mu\nu}V^{A,\mu\nu}$$
$$+ i\bar{f}_L\gamma^\mu \mathcal{D}_\mu f_L + i\bar{f}_R\gamma^\mu \mathcal{D}_\mu f_R \tag{79}$$
$$+ \mathcal{L}'_{dim>4}\left(W^i_{\mu\nu}, V^A_{\mu\nu}; f_L, \mathcal{D}_\mu f_L; f_R, \mathcal{D}_\mu f_R\right)$$

where summation over $i = 1,2,3$, $a = 1,\ldots,8$, $A = 1,\ldots,15$ and the fermions is implied. W_μ^i and $W_{\mu\nu}^i$ are as defined in Eqs. (16) and (20), respectively. In addition, we now have for $\alpha, \alpha_c \neq 0$ the substitutions

$$\begin{aligned} V_\mu^{15} &= v_\mu^{15} + \frac{e}{g_V}\sqrt{\frac{2}{3}}\,A_\mu \\ V_\mu^a &= v_\mu^a + \frac{g_c}{g_V}G_\mu^a \quad , \quad a = 1,\ldots,8 \\ V_\mu^{9,\ldots,14} &= v_\mu^{9,\ldots,14} . \end{aligned} \tag{80}$$

$V_{\mu\nu}^A$ denote the usual SU(4)-Yang-Mills field-strength tensors, involving the SU(4) effective coupling g_V and structure constants C^{ABC}

$$V_{\mu\nu}^A = \partial_\mu V_\nu^A - \partial_\nu V_\mu^A - g_V C^{ABC} V_\mu^B V_\nu^C . \tag{81}$$

The covariant derivatives read

$$\begin{aligned} \mathcal{D}_\mu f_L &= \left(\partial_\mu + i g_W W_\mu^i \frac{\tau^i}{2} + i g_V V_\mu^A \frac{\lambda^A}{2}\right) f_L \\ \mathcal{D}_\mu f_R &= \left(\partial_\mu + i g_c G_\mu^a \frac{\lambda^a}{2} + i e A_\mu y_R\right) f_R . \end{aligned} \tag{82}$$

Again, the γ-w^3 and γ-y "current-mixing" terms can be removed by the same linear field transformation as in Sect 5 (given in Ref.14). The gluon-v_8 "current-mixing" terms

$$-\frac{1}{4}\frac{g_c}{g_V}\left(v_{\mu\nu}^a G^{a,\mu\nu} + G_{\mu\nu}^a v^{a,\mu\nu}\right) \tag{83}$$

contained in Eq. (79), are removed by a linear transformation analogous to Eq. (43a)

$$\begin{aligned} G_\mu^a &= G_\mu^{a,diag} - tg\Theta_c\, v_\mu^{a,diag} \\ v_\mu^a &= v_\mu^{a,diag}/\cos\Theta_c \end{aligned} \tag{84}$$

with
$$\sin\Theta_c = \frac{g_c}{g_V} < 1 . \tag{85}$$

In this <u>diagonal</u> basis one then finds from Eqs. (79,82,84) the following couplings of gluons and color octet vector bosons to quarks

$$\begin{aligned} \mathcal{L}_{q\bar{q}} = &- g_c G_\mu^{a,diag}\left[\bar{q}_L\gamma^\mu\frac{\lambda^a}{2}q_L + \bar{q}_R\gamma^\mu\frac{\lambda^a}{2}q_R\right] \\ &- g_c v_\mu^{a,diag}\left[ctg\Theta_c\,\bar{q}_L\gamma^\mu\frac{\lambda^a}{2}q_L - tg\Theta_c\,\bar{q}_R\gamma^\mu\frac{\lambda^a}{2}q_R\right] \\ &+ \ldots \end{aligned} \tag{86}$$

The gluons couple in the correct, vectorlike manner. The coupling of the color-octet bosons v_8^{diag} is reminiscent of the coupling of the Z boson (with e, ϑ_W replaced by g_c, Θ_c).

In terms of the four parameters m_W, g_W and m_V, g_V (and the two known couplings e, g_c) a well-defined mass hierarchy emerges within each of the multiplets $\vec{w}_\mu$ and $\vec{v}_\mu$. The <u>lightest</u> vector bosons of a given multiplet are those which have <u>no mixing</u> with photons or gluons, i.e. in the $SU(2)_{WI}$ sector the charged W's

$$m(W^\pm) = m_W \tag{87}$$

and in the $SU(4)^L$ sector the color triplets

$$m(v_3) = m_V \tag{88}$$

The color octet bosons v_8, after diagonalization, have a higher mass due to mixing with gluons (cf. Eq. (84))

$$m(v_8^{diag}) = m_V/\cos\Theta_c > m(v_3) . \tag{89}$$

Since the $SU(4)^L$ and $SU(2)_{WI}$ sectors are linked through the electromagnetic charge operator, the mass formula for the Z and y^{diag} bosons is more complicated; it is given in Eq. (55) with $g_Y = \sqrt{\frac{3}{2}}\, g_V$.

It may be instructive to consider the (plausible) special case

$$\left(\frac{e}{g_Y}\right)^2 \ll \frac{1-\left(\frac{e}{g_W}\right)^2}{1+\left(\frac{e}{g_W}\right)^2} \sim 0.63. \qquad (90)$$

The square root in Eq. (55) may then be approximated such that

$$\frac{m_Z}{m_W} \simeq \left(\frac{1-\left(\frac{e}{g_Y}\right)^2}{1-\left(\frac{e}{g_W}\right)^2-\left(\frac{e}{g_Y}\right)^2}\right)^{1/2} \sim \frac{1}{\cos\Theta_W}, \qquad \frac{m(y^{diag})}{m(v_3)} = \left(\frac{1-\left(\frac{e}{g_W}\right)^2}{1-\left(\frac{e}{g_W}\right)^2-\left(\frac{e}{g_Y}\right)^2}\right)^{1/2} \sim 1. \qquad (91)$$

In this case the y^{diag} and v_3 vector bosons are almost mass degenerate while the color octet bosons are substantially heavier.

Complete information about the four parameters m_W, g_W and m_V, g_V can, in principle, come from precision measurements of G_F, $\sin^2\Theta_W$, m_W and m_Z. In order to identify G_F and $\sin^2\Theta_W$ as functions of the parameters m_W, g_W and m_V, g_V, the low-energy four-fermion limit of $\mathcal{L}_{eff}^{\alpha,\alpha_c \neq 0}$ has to be taken. Let us only indicate the procedure qualitatively here. The details will be presented elsewhere.

Along the lines of Ref. 14 one finds after an appropriate Fierz transformation

$$\mathcal{L}_{eff}^{weak} \xrightarrow[\Lambda_c \ll E \ll m_{W,Z}]{} -\frac{1}{2}\, G_F^{exp}\,(\text{charged current int.} + \text{neutral current int.}(\sin^2\Theta_W^{exp})) + \text{scalar currents} \qquad (92)$$

This gives the desired relations of G_F^{exp} and $\sin^2\Theta_W^{exp}$ in terms of m_W, g_W and m_V, g_V. Measurements of m_W and m_Z together with Eq. (55) complete the system of equations.

Given the good agreement of the GSW model with the data so far available, one expects the masses of the new composite vector bosons to lie at least in the few hundred GeV range. See also Ref. 29 and the footnote on page 21 in this context.

The dominant decay modes of the new vector bosons are

$$\begin{array}{lcl} y^{diag} & \longrightarrow & q\bar{q},\ \ell\bar{\ell} \\ v_3 & \longrightarrow & q\bar{\ell} \\ v_8^{diag} & \longrightarrow & q\bar{q},\ \text{gluons} \end{array} \qquad (93)$$

with couplings to be inferred from $\mathcal{L}_{eff}$, Eqs. (78,79,86). The bosons y^{diag} and v_8^{diag} are easily produced in $\bar{p}p$ collisions. For the production of the bosons v_3 a high energy ep machine (HERA?) is more suited. The most exotic decay modes are those of the lightest vector bosons v_3, into a jet and a lepton which could be a neutrino. In the latter case the signature is a monojet event with high missing energy.

7. SUMMARY AND CONCLUSIONS

The present work is based on the idea that - unlike the standard GSW model - quarks, leptons and the W,Z vector bosons are all composite. Their weak interactions are viewed, in analogy to strong interactions among composite hadrons, as residual interactions (of range $\frac{1}{\Lambda} \sim G_F^{1/2}$) caused by an underlying confining hypercolor gauge theory for preons.

The aim of the present investigation was to set up a systematic effective Lagrangian approach for the weak interactions in the "low-energy" regime, $E \lesssim \Lambda \sim G_F^{-1/2} \sim 300$ GeV, where the composites can be described in terms of local fields.

Unfortunately, a direct link between the underlying hypercolor gauge theory and the corresponding $\mathcal{L}_{eff}$ cannot be established at present.

Our guidelines for constructing $\mathcal{L}_{eff}$ came from two sources

i) $\mathcal{L}_{eff}$ is to incorporate a maximum of information which at present may be extracted from specific preon gauge models.

ii) We generalized a "low-energy" concept which has already once led to successfully determine $\mathcal{L}_{eff}$ in the prototype case of low-energy strong interactions.

This led us to the following input requirements from which all our results were derived.

(i) We abstracted from preon models the specification of the global chiral symmetry G of weak interactions for α and $\alpha_{color} \to 0$ and the classification of the massless composite ground state fermions (quarks and leptons) with respect to G. Furthermore we required the (plausible) existence of massive composite vector bosons in the adjoint rep. of G, all of which can be related to hypercolor singlet composite operators in terms of preon fields. This establishes the maximal link to preon models incorporating 't Hooft's anomaly constraints. For $\alpha, \alpha_{color} \neq 0$, we required local gauge invariance with respect to the $U(1)_{em} \times SU(3)_c$ subgroup of G.

(ii) The powerful requirement of current-field identities for the exact local symmetry currents. This implies the conventional vector boson dominance in the photon sector as well as generalization to the gluon sector involving the proportionality of color octet currents and color octet vector boson fields.

As a result of (i) and (ii) the effective Lagrangian is fixed in terms of a few parameters. It is of the massive Yang-Mills type with respect to the global group G and involves two parameters, a mass and a "quasi-gauge" coupling, for each simple factor in G.

Two applications were studied in more detail. A restriction to $G = SU(2)_{WI}$ and the known particles q, l, $W^{\pm}$, Z led to an $\mathcal{L}_{eff}$ which closely mimics GSW predictions in the absence of a physical Higgs scalar. Instrumental for this strong conclusion were two properties i) that G is a chiral symmetry, keeping the quarks and leptons massless on the scale m_W, and ii) that the effective $SU(2)_{WI}$-"quasi-gauge" coupling is small.

The second application specialized to the symmetry $G = SU(2)_{WI} \times SU(4)_L$ as abstracted from a popular class of Abbott-Farhi type models. It involves the familiar $SU(2)_{WI}$-triplet of (composite) w bosons and a new $SU(4)_L$-15 plet of composite vector bosons with the color decomposition $\underline{15} = \underline{8} + \underline{3} + \underline{\bar{3}} + \underline{1}$ $(= v_8 + v_3 + v_{\bar{3}} + y)$. Vector boson dominance, relating the electromagnetic current to the w^3 and y fields and the color octet currents to the v_8 fields, was implemented. The resulting $\mathcal{L}_{eff}$ involves besides the small, known gauge couplings α and $\alpha_c(G_F^{-1/2})$ as free parameters two masses and two "quasi-gauge" couplings m_W, m_V and g_W, g_V. A hierarchy of mass-splitting among the two vector boson multiplets, arising an account of $U(1)_{em} \times SU(3)_c$ gauge interactions, was derived. The couplings of the new vector bosons to quarks and leptons were determined. Implications for high precision experiments for $E \lesssim m_{W,Z}$ as well as $\bar{p}p$ collider and ep (HERA) experiments were briefly discussed.

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FIGURE CAPTIONS

Fig. 1a: Mixing of the composite vector boson w^3 with the photon as resulting from the current-field identity for the weak isovector part of the electromagnetic current.

Fig. 1b: Mixing of composite, color octet vector bosons v_8 with gluons as resulting from the current-field identity for the color octet current.

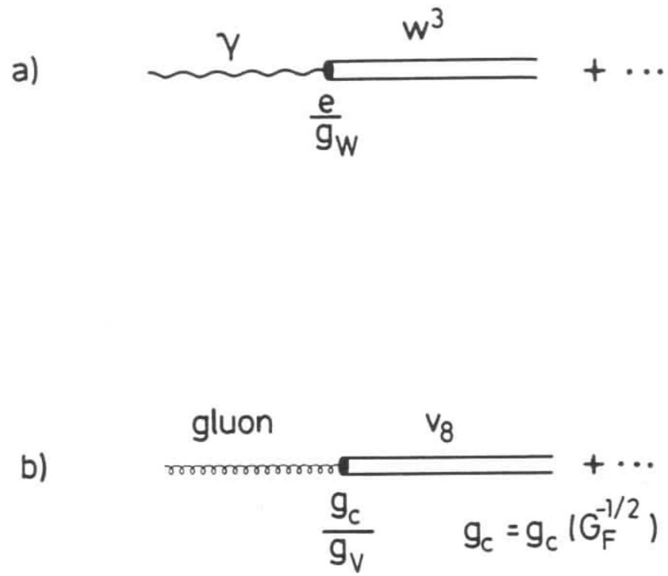

Fig. 1