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SCALAR CONDENSATES AND PHYSICAL SPECTRUM

IN A SUPERSYMMETRIC GAUGE MODEL

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ABSTRACT

We investigate a supersymmetric model with $SU(2)$ gauge interaction and global $U(6)$ invariance which has recently been proposed as a toy model for composite quarks, leptons and intermediate vector bosons. The vacuum structure of the model is strongly restricted. We evaluate the dominant contributions of scalar condensates to the two-point functions of the $U(6)$ currents and study the corresponding SVZ sum rules for the physical spectrum of the theory. Our results suggest that the $U(6)$ symmetry is spontaneously broken to $SU(2) \times U(4)$. The simplest Ansatz for the physical spectrum, which is compatible with the sum rules, contains only the states present in the complementary (Higgs) picture. Implications of our results for preon models with composite W -bosons are briefly discussed.

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1. INTRODUCTION

In preon models of quarks and leptons¹⁾ it is natural to consider the possibility that also the intermediate vector bosons are composite. It is known that composite W- and Z-bosons are difficult to distinguish from elementary gauge particles if their masses are small compared to the new substructure scale: a dynamical γ -W³ mixing²⁾ can lead to the observed neutral current structure and vector meson dominance³⁾ implies the standard model mass predictions. It is unknown, however, how one can obtain dynamically composite W-bosons which are light compared to the other QCD-like vector bosons which are usually predicted in preon models.

The central question of quark-lepton substructure is the dynamical origin of (almost) massless composite fermions, a problem whose solution may involve supersymmetry¹⁾. In Ref. 4) a model has been constructed in which the left-handed quarks and leptons of one family are quasi-Goldstone fermions, i.e., superpartners of Goldstone bosons which arise from the spontaneous breaking of a global U(6) symmetry to the subgroup SU(2) $\times$ U(4). The structure of the low energy effective Lagrangian of the quasi-Goldstone fermions led to the conjecture⁴⁾ that among the expected 36 composite vector supermultiplets the SU(2) triplet may indeed be the lightest due to the effect of the scalar condensates which cause the symmetry breaking.

A means to study the influence of the vacuum structure on vector boson masses is the SVZ sum rule technique⁵⁾. Most of the applications to preon models⁶⁻⁸⁾ have treated the bare loop as dominant contribution, as it is the case in QCD. More recently it has been realized, however, that in theories with scalar condensates, like the model of Abbott and Farhi⁹⁾, power "corrections" may be dominant^{10),11)}. In this case a picture of composite vector bosons emerges which is substantially different from that of the ρ -meson in QCD.

In this paper we follow the approach of Ref. 10) to study the vector supermultiplets which occur in the U(6) model of Ref. 4). In Section 2 we review the main features of the model and in Section 3 we discuss in some detail its vacuum structure, making use of results obtained for general SQCD theories. Section 4 deals with the two-point functions of the U(6) currents and the associated SVZ sum rules. In Section 5 we summarize our results and draw some conclusions concerning preon models with composite W-bosons.

2. THE U(6) MODEL

A supersymmetric preon model with residual weak interactions, which is based on the group $U(6)^{*)}$, has been studied in Ref. 4). The spontaneous breaking of $U(6) \rightarrow SU(2) \times U(4)$ produces 17 Goldstone bosons which can be accommodated in nine chiral supermultiplets ϕ_{ia} ($i=1,2$; $a=3\dots 6$) and ϕ , transforming as $(2^*, 4^*)$ and $(1, 1)$ with respect to the unbroken subgroup. The four doublets ϕ_{ia} contain the colour triplet of quark doublets and the lepton doublet; the singlet ϕ corresponds to an additional state, the "novino". The $SU(2)$ and $SU(4)$ factors in the unbroken subgroup are identified as weak isospin and the Pati-Salam symmetry¹²⁾, which guarantees the universality of the residual weak interactions.

The low energy interactions of Goldstone bosons and quasi-Goldstone fermions are described by an effective Lagrangian which is invariant under the nonlinearly realized $U(6)$ symmetry. For the fermionic part one obtains⁴⁾:

$$\mathcal{L}_{\text{eff}}^{(q,l)} = \frac{1}{2f_1^2} \left(\bar{\psi}_L^{ia} \gamma^\mu \left(\frac{1}{i} \vec{\tau} \right)_i^j \psi_{Lja} \right)^2 + \frac{f_2^2 - f_1^2}{2f_2^2} \left(\frac{1}{i} \bar{\psi}_L^{ia} \gamma^\mu \psi_{Lia} \right)^2 + \dots, \quad (1)$$

where the scales f_1 and f_2 are the "decay constants" of the Goldstone multiplets ϕ and ϕ_{ia} . In the limit of vanishing electromagnetic coupling the weak interactions of the standard model are an isovector current-current interaction which corresponds to the first term in Eq. (1). It is remarkable that the unwanted isoscalar term in Eq. (1) is indeed suppressed in the case $f_1 \approx f_2$, which one expects if the symmetry breaking occurs in a confining preon theory with one hypercolour scale.

Such a preon model is easily constructed. A global $U(6)$ symmetry is obtained by introducing six chiral multiplets χ_α^p ($\alpha = 1\dots 6$), where p denotes the hypercolour index. The simplest choice of a confining, anomaly-free gauge interaction corresponds to the hypercolour group $SU(2)_{\text{HC}}$ under which the chiral multiplets χ_α^p transform as doublets, i.e., $p = 1, 2$. Due to supersymmetry the full global symmetry of the Lagrangian contains an additional R symmetry,

*) The group $U(6)$ corresponds to the case of one family. An extension to the group $U(4n_F+2)$ for n_F families is straightforward.

$$G_{cl} = U(6) \times U(1)_R \quad . \quad (2)$$

Instanton effects reduce this invariance of the Lagrangian to

$$G_{qu} = SU(6) \times U(1)_X \quad , \quad (3)$$

where X is given by

$$X = T^a_\alpha + 3 R \quad . \quad (4)$$

R and T^a_α are the generators of $U(1)_R$ and the $U(1)$ factor in $U(6)$.

The simplest composite hypercolour singlet operators are the chiral superfields

$$\Phi_{\alpha\beta} = \varepsilon_{pq} \chi^p_\alpha \chi^q_\beta \quad , \quad (5)$$

and the vector superfields

$$J^a_\beta = \bar{\chi}^a_p (e^{2gV})^p_q \chi^q_\beta \quad , \quad (6)$$

where $V = \frac{1}{2} \tau^I V^I$ denotes the $SU(2)_{HC}$ gauge fields. The vector multiplets J^α_β are the conserved currents^{*)} of the theory (cf. Ref.13)), i.e.,

$$D^2 J^a_\beta = \bar{D}^2 J^a_\beta = 0 \quad . \quad (7)$$

Eq. (7) contains the familiar current conservation

$$\partial^\mu j^a_{\beta\mu} = 0 \quad , \quad (8)$$

with

$$j^a_{\beta\mu} = \bar{\eta}^a_p \overleftrightarrow{\partial}_\mu \eta^p_\beta + i \tilde{\chi}^{*a}_p \overleftrightarrow{\partial}_\mu \tilde{\chi}^p_\beta + g V^I_\mu \tilde{\chi}^{*a}_p (T^I)^p_q \tilde{\chi}^q_\beta \quad , \quad (9)$$

where $\tilde{\chi}^p_\alpha$ and η^p_α denote the scalars and Weyl fermions which are contained in the

*) We ignore here complications due to the anomaly of the $U(1)$ factor.

chiral multiplets χ_α^p . Vacuum expectation values of $\Phi_{\alpha\beta}$ and J_β^α can break the $U(6)$ invariance to various subgroups such as $SU(2) \times U(4)$, $Sp(6)$ or $(SU(2))^3$. A remarkable feature of the model is that the case of unbroken $SU(6) \times U(1)_X$ symmetry is also consistent with 't Hooft's anomaly conditions¹⁴⁾. The triangle anomalies of the $SU(6) \times U(1)_X$ currents at the preon level are matched at the bound state level by the 15 fermions contained in $\Phi_{\alpha\beta}$ ¹⁵⁾.

As our model contains scalar fields in the fundamental representation of the gauge group we can study some of its properties in the "complementary picture"^{14),16)} by giving vacuum expectation values to the fields χ_α^p . This "Higgs phase" is nothing but the supersymmetric standard model without right-handed fields, so that our $U(6)$ model may be regarded as a supersymmetric version of the Abbott-Farhi model⁹⁾. Unbroken supersymmetry requires a vanishing D-term,

$$\langle \chi_\rho^{*d} \rangle (\tau^I)_\rho^q \langle \chi_a^q \rangle = 0, \quad I=1 \dots 3. \quad (10)$$

This constraint leads to $SU(2) \times U(4)$ as the only possible global symmetry in the "Higgs phase".

As we will see in the next section, the analysis of the "confining phase" yields the same result for the unbroken global symmetry, in agreement with the idea of "complementarity"¹⁶⁾.

3. THE VACUUM STRUCTURE

The model described in the previous section corresponds to SQCD with two colours and three flavours. As the fundamental representation of $SU(2)$ is pseudoreal, the global symmetry is $U(6)$ rather than $U(3) \times U(3)$. The composite fields

$$\Phi_{\alpha\beta} = \varepsilon_{pq} \chi_\alpha^p \chi_\beta^q, \quad J_\beta^d = \bar{\chi}_\rho^d (e^{2gV})_\rho^q \chi_\beta^q,$$

can have the vacuum expectation values ($i=1,2$; $a_1=3,4$; $a_2=5,6$)

$$\begin{aligned}
 \langle \bar{\Phi}_{\bar{i}\bar{j}} \rangle &= V_1^2 \varepsilon_{\bar{i}\bar{j}} \quad , \quad \langle \bar{\Phi}_{a,b_1} \rangle = V_2^2 \varepsilon_{a,b_1} \quad , \\
 \langle \bar{\Phi}_{a_2,b_2} \rangle &= V_3^2 \varepsilon_{a_2,b_2} \quad , \quad \langle J_{\bar{j}}^i \rangle = \tilde{V}_1^2 \delta_{\bar{j}}^i \quad , \\
 \langle J_{b_1}^{a_1} \rangle &= \tilde{V}_2^2 \delta_{b_1}^{a_1} \quad , \quad \langle J_{b_2}^{a_2} \rangle = \tilde{V}_3^2 \delta_{b_2}^{a_2} \quad .
 \end{aligned} \tag{11}$$

where we have assumed that the vector condensates, like the chiral condensates, break the $U(6)$ symmetry to the subgroup $(SU(2))^3$. A further bilinear condensate can be formed by the gauginos,

$$\langle \lambda^{\bar{I}} \lambda^{\bar{I}} \rangle = - \langle W^{\bar{I}d} W_{\bar{d}}^{\bar{I}} \rangle = V^3 \quad , \tag{12}$$

where $W_{\alpha}^{\bar{I}}$ is the chiral multiplet¹³⁾

$$W_{\bar{d}}^{\bar{I}} = \text{tr} \left[\tau^{\bar{I}} \left(-\frac{1}{4} \right) \bar{D}^2 e^{-V} D_{\bar{d}} e^V \right] \quad . \tag{13}$$

Information on the chiral condensates can be obtained by breaking the $U(6)$ invariance explicitly to $(SU(2))^3$ through the superpotential

$$g_m = -\frac{1}{2} m_1 \varepsilon^{\bar{i}\bar{j}} \bar{\Phi}_{\bar{i}\bar{j}} - \frac{1}{2} m_2 \varepsilon^{a,b_1} \bar{\Phi}_{a,b_1} - \frac{1}{2} m_3 \varepsilon^{a_2,b_2} \bar{\Phi}_{a_2,b_2} \tag{14}$$

and considering the chiral limit $m_1 \rightarrow 0$. The effective Lagrangian approach¹⁷⁾ implies for the chiral condensates the scaling behaviour

$$\begin{aligned}
 V_{\bar{i}}^2 &\sim \frac{1}{m_{\bar{i}}} (m_1 m_2 m_3)^{\frac{1}{2}} \quad , \quad \bar{i} = 1 \dots 3 \quad , \\
 V^3 &\sim (m_1 m_2 m_3)^{\frac{1}{2}} \quad .
 \end{aligned} \tag{15}$$

These results have been further supported by analyticity arguments¹⁸⁾ and SUSY instanton calculations¹⁹⁾. If we demand that in the chiral limit all vacuum expectation values remain finite, Eq. (15) leads to two different cases for vanishing m_1 :

$$(i) \quad V_1 = V_2 = V_3 = 0, \quad (16)$$

$$(ii) \quad V_1 \neq 0, \quad V_2 = V_3 = 0.$$

The first case, which corresponds to unbroken $U(6)$ symmetry, results from a uniform limit, i.e., $m_1 \sim m_2 \sim m_3 \rightarrow 0$; the second case, which corresponds to the breaking $U(6) \rightarrow SU(2) \times U(4)$, is obtained from the non-uniform limit $m_1 \sim m_2 m_3 \rightarrow 0$.

With respect to the vector condensates, it is useful to consider the SUSY Dashen formulae^{20),21)}

$$f_A (M^2)_{AB} f_B = \frac{1}{2} \langle [X_A, [X_B, \mathcal{L}_F]] \rangle, \quad (17a)$$

$$f_A M_{AB} f_B = \langle [X_A, [X_B, \mathcal{L}_S]] \rangle; \quad (17b)$$

$\mathcal{L}_F(\mathcal{L}_S)$ are the fermionic (scalar) symmetry breaking parts of the superpotential, M_{AB} the pseudo-Goldstone boson mass matrix, X_A the broken charges and f_A the related decay constants. X_A can be expressed in terms of the current $J_A^{21)$,

$$X_A = \int d^3x \, J_A^0, \quad (18a)$$

$$J_A^\mu = -\frac{1}{4} \sigma_{\dot{a}a}^\mu [\bar{D}^{\dot{a}} \bar{D}^a] J_A, \quad (18b)$$

which is connected with the Goldstone chiral multiplet ϕ_A through the current field identity²¹⁾

$$-\frac{1}{4} \bar{D}^2 J_A = M_A f_A \phi_A. \quad (19)$$

Let us apply the SUSY Dashen formulae to our model. The symmetry breaking parts of the Lagrangian read

$$\mathcal{L}_S = \frac{1}{2} \frac{\partial^2 g_m}{\partial \chi_\alpha^p \partial \chi_\beta^q} \tilde{\chi}_\alpha^p \tilde{\chi}_\beta^q, \quad (20a)$$

$$\mathcal{L}_F = \frac{1}{2} \frac{\partial^2 g_m}{\partial \chi_\alpha^p \partial \chi_\beta^q} \eta_\alpha^p \eta_\beta^q + \text{c.c.}, \quad (20b)$$

where $\tilde{\chi}_\alpha^p$ (η_α^p) is the scalar (fermionic) component of the chiral multiplet χ_α^p . For the currents J_i^1 , $J_{a_1}^1$, $J_{a_2}^1$, $J_i^{a_1}$ and $J_i^{a_2}$, which correspond to the broken $U(6)$ generators $\sqrt{2}X$, $X_{a_1}^1$, $X_{a_2}^1$, $X_i^{a_1}$ and $X_i^{a_2}$ (cf. Ref. 4)), one has:

$$-\frac{1}{4} \bar{D}^2 J_i^1 = -\sqrt{2} M_1 f_1 \Phi, \quad (21a)$$

$$-\frac{1}{4} \bar{D}^2 J_{a_1}^1 = \varepsilon^{ij} M_2 f_2 \Phi_{ja_1}, \quad (21b)$$

$$-\frac{1}{4} \bar{D}^2 J_{a_2}^1 = \varepsilon^{ij} M_3 f_3 \Phi_{ja_2}. \quad (21c)$$

The Dashen formula (17a) involves a bilinear fermionic condensate. As the vacuum expectation values of the F-components of $\chi_\alpha^p \chi_\beta^q$ vanish, the fermionic condensates can be replaced by scalar condensates, for instance,

$$\langle \eta_i^p \eta_{a_1}^q \rangle = \varepsilon^{pr} \varepsilon_{ik} w_1 \langle \tilde{\chi}_r^{*k} \tilde{\chi}_{a_1}^q \rangle + \varepsilon^{qr} \varepsilon_{a_1 b_1} w_2 \langle \tilde{\chi}_i^p \tilde{\chi}_r^{*b_1} \rangle. \quad (22)$$

From Eqs. (11), (17), (20), (21) and (22), one obtains the following two sets of equations (we have chosen all parameters to be real):

$$\begin{aligned} M_1 f_1^2 &= 2 w_1 v_1^2, \\ M_2 f_2^2 &= w_1 v_1^2 = w_2 v_2^2, \\ M_3 f_3^2 &= w_1 v_1^2 = w_3 v_3^2, \end{aligned} \quad (23a)$$

and

$$\begin{aligned}
 M_1^2 f_1^2 &= 4 m_1^2 \tilde{V}_1^2, \\
 M_2^2 f_2^2 &= m_1^2 \tilde{V}_1^2 + m_2^2 \tilde{V}_2^2, \\
 M_3^2 f_3^2 &= m_1^2 \tilde{V}_1^2 + m_3^2 \tilde{V}_3^2.
 \end{aligned}
 \tag{23b}$$

Let us first consider the non-uniform chiral limit $m_1 \sim m_2 m_3 \rightarrow 0$ which leads to the $SU(2) \times U(4)$ case (cf. Eq. (16)): $v_1 \rightarrow C \neq 0$ and $v_2, v_3 \rightarrow 0$. In this limit one expects $f_1 \neq 0$, $f_2 = f_3 = 0$ and therefore

$$\begin{aligned}
 \tilde{V}_1^2 &= \frac{v_1^4}{f_1^2} \rightarrow \tilde{C}_1 \neq 0, \\
 \tilde{V}_2^2 &\leq \tilde{C}_2 \sim v_2^4 \rightarrow 0, \\
 \tilde{V}_3^2 &\leq \tilde{C}_3 \sim v_3^4 \rightarrow 0,
 \end{aligned}
 \tag{24}$$

i.e., the vector condensate $\tilde{V}_1^2$ ($\tilde{v}_{2,3}^2$) is proportional to (bounded by) the square of the chiral condensate v_1^2 ($v_{2,3}^2$) in the chiral limit^{*}). Therefore, for vanishing m_1 , also the vector condensates break the $U(6)$ symmetry to the subgroup $SU(2) \times U(4)$.

In the uniform chiral limit $m_1 \sim m_2 \sim m_3 \rightarrow 0$, all chiral condensates vanish, i.e., $v_1 \sim v_2 \sim v_3 \rightarrow 0$, and there is no reason anymore to expect that the decay constants f remain different from zero. Although in this case it is not possible to draw a firm conclusion from Eqs. (23), it seems most natural that all vector condensates vanish, i.e., $\tilde{V}_1 \sim \tilde{V}_2 \sim \tilde{V}_3 \rightarrow 0$. This is a consequence of a vanishing D-term (cf. Eqs. (26), (27)) and the factorization hypothesis for gauge invariant operators^{**}):

*) For SQCD with less flavours than colours, constraints on vector condensates have recently been obtained by Konishi²²⁾.

**) Whenever possible without causing confusion we use the same symbol for a chiral superfield and its scalar component.

$$\begin{aligned}
 \langle \vec{D}^2 \rangle &= \langle (\chi_p^{*d} (\frac{1}{i} \vec{\tau})_q^p \chi_d^q) (\chi_r^{*b} (\frac{1}{i} \vec{\tau})_s^r \chi_b^s) \rangle \\
 &= \frac{1}{2} \langle (\chi_p^{*d} \chi_p^p) (\chi_q^{*b} \chi_q^b) - \frac{1}{2} (\chi_p^{*d} \chi_d^p)^2 \rangle \\
 &= -2 (\tilde{V}_1^2 \tilde{V}_2^2 + \tilde{V}_1^2 \tilde{V}_3^2 + \tilde{V}_2^2 \tilde{V}_3^2)
 \end{aligned} \tag{25}$$

We thus expect that in the case of unbroken $U(6)$ symmetry all scalar condensates vanish. From Eqs. (15), (22) and (24) it is apparent that in the chiral limit all bilinear fermionic condensates are zero.

Our analysis of the vector multiplets in the following section will involve higher dimensional condensates. A constraint on quartic scalar condensates follows from

$$\begin{aligned}
 &\frac{1}{4} \langle W^{Ia} W_a^I |_{\theta\theta} + \bar{W}_a^I \bar{W}^{Ia} |_{\bar{\theta}\bar{\theta}} \rangle \\
 &= \langle -\frac{1}{4} F_{\mu\nu}^I F^{I\mu\nu} - i \lambda^I \sigma^\mu \partial_\mu \bar{\lambda}^I + \frac{1}{2} D^I D^I \rangle = 0
 \end{aligned} \tag{26}$$

As the gaugino condensate $\langle \lambda^I \lambda^I \rangle$ vanishes for $m_i=0$ (cf. Eq. (15)), it appears reasonable to neglect possible contributions from $\langle \lambda^I \bar{\lambda}^I \rangle$ and, correspondingly, $\langle F_{\mu\nu}^I F^{I\mu\nu} \rangle$. Unbroken supersymmetry then requires

$$\langle \vec{D}^2 \rangle = \langle (\chi_p^{*d} (\frac{1}{i} \vec{\tau})_q^p \chi_d^q)^2 \rangle = 0, \tag{27}$$

which, for classical fields, is the gauge invariant version of Eq. (10).

We have shown that the scalar fields χ_a^p ($a=3\dots 6$) do not contribute to bilinear condensates. We will assume that also higher dimensional condensates involving the fields χ_a^p vanish and that condensates of the scalar fields χ_i^p ($i=1,2$) are flavour $SU(2)$ symmetric. This implies (we drop hypercolour indices):

$$\langle (\chi^{*i} \vec{\tau} \chi_j) (\chi^{*k} \vec{\tau} \chi_k) \rangle = 2 \delta_j^i \langle \vec{D}^2 \rangle = 0, \tag{28}$$

and

$$\begin{aligned} & \langle (\chi^{*i} \bar{\chi}_j) (\chi^{*k} \bar{\chi}_l) \rangle \\ & \equiv \frac{1}{2} (\delta_l^i \delta_j^k - \frac{1}{2} \delta_j^i \delta_l^k) \langle (\chi^{*} \bar{\chi})^2 \rangle \end{aligned} \quad (29)$$

We will assume condensates of six scalars χ_1^p to be zero if they contain a gauge invariant factor whose vacuum expectation value vanishes. Thus we have $\langle \chi^{*} \chi \equiv \sum_l \chi^{*l} \chi_l \rangle$

$$\langle (\chi^{*i} \bar{\chi}_j) (\chi^{*k} \bar{\chi}_l) (\chi^{*} \chi) \rangle = 2 \delta_j^i \langle \bar{\chi}^l \rangle \langle \chi^{*} \chi \rangle = 0, \quad (30)$$

$$\langle (\chi^{*i} \bar{\chi}_j) (\chi^{*u} \bar{\chi}_w) (\chi^{*k} \chi_l) \rangle = 4 \delta_j^i \delta_l^k \langle \bar{\chi}^u \rangle \langle \chi^{*} \chi \rangle = 0, \quad (31)$$

and

$$\begin{aligned} & \langle (\chi^{*i} \bar{\chi}_j) (\chi^{*k} \bar{\chi}_l) (\chi^{*} \chi) \rangle \\ & \equiv \frac{1}{2} (\delta_l^i \delta_j^k - \frac{1}{2} \delta_j^i \delta_l^k) \langle (\chi^{*} \bar{\chi})^2 (\chi^{*} \chi) \rangle \end{aligned} \quad (32)$$

We note that in Eqs. (30), (31) and (32) only the isovector channel is non zero.

If one wants to express higher dimensional condensates in terms of bilinear ones, further assumptions are necessary (cf. Refs. 5), 10)). For instance, factorization of vacuum expectation values of products of gauge invariant operators would imply

$$\begin{aligned} \langle \bar{\chi}^i \chi^j \rangle &= \langle (\chi^{*i} \bar{\chi}_p (\frac{1}{2} \bar{\chi})^p \chi^q \chi_i^q)^2 \rangle \\ &= \frac{1}{4} \langle (\chi^{*i} \chi_i^p)^2 - 1 \varepsilon^{ij} \varepsilon_{pq} \chi_i^p \chi_j^q \rangle \\ &= \tilde{V}_1^4 - V_1^4, \end{aligned} \quad (33)$$

and therefore (cf. Refs. 24), 27))*)

*) Based on duality, Narison²³⁾ has argued in favour of the validity of Eq. (34).

$$\tilde{V}_1^2 = V_1^2 = f_1^2 \quad (34)$$

4. COMPOSITE VECTOR BOSONS

Given the scalar condensates discussed in the previous section, we can now study their influence on the physical spectrum of the theory by means of SVZ sum rules^{5),10)}. In the standard manner, the two-point functions of the currents J_β^α in the Euclidean region can be evaluated using dispersion relations and an Ansatz for the physical spectrum on the one side, and in perturbation theory on the other side. The structure of the vacuum is incorporated in the perturbative expansion through condensates. Duality relations can then be used to compute some properties of the physical spectrum.

The calculation of the two-point functions of the currents J_β^α is straightforward. In the Wess-Zumino gauge, the currents and the interaction Lagrangian read ($V = \frac{1}{2}\tau^{II}V^I$)

$$\begin{aligned} J_\beta^a &= \bar{\chi}_\rho^a (e^{2gV})_\rho^q \chi_\beta^q \\ &= \bar{\chi}_\rho^a \chi_\beta^p \left[1 + \frac{g^2}{2}(V^I)^2 \right] + g \bar{\chi}_\rho^a (\tau^I)_\rho^q \chi_\beta^q V^I, \end{aligned} \quad (35)$$

and

$$\mathcal{L}_I = \int d^4\theta \left[g \bar{\chi}_\rho^a (\tau^I)_\rho^q \chi_\rho^q V^I + \frac{g^2}{2} \bar{\chi}_\rho^a \chi_\rho^p (V^I)^2 \right] \quad (36)$$

The propagators of chiral and vector superfields are given by¹³⁾

$$\langle T(\chi_a^P(1) \bar{\chi}_q^P(2)) \rangle = i \delta_a^P \delta_q^P P_i \delta(1,2) , \quad (37a)$$

$$\langle T(V^A(1) V^B(2)) \rangle = \frac{i}{2} \delta^{AB} \frac{-1}{\square} P_T \delta(1,2) , \quad (37b)$$

with $\delta(1,2) = \delta^4(x_1 - x_2) \delta^4(\theta_1 - \theta_2) , \quad \delta^4(\theta_1 - \theta_2) = \delta^2(\theta_1 - \theta_2) \delta^2(\bar{\theta}_1 - \bar{\theta}_2) ,$

$$P_i = \frac{1}{16} \frac{\bar{D}^2 D^2}{\square} , \quad P_2 = \frac{1}{16} \frac{D^2 \bar{D}^2}{\square} , \quad P_T = -\frac{1}{8D} D \bar{D}^2 D ; \quad (37c)$$

in (37b) we have chosen the Landau gauge.

The calculation of the time-ordered product of the currents J_β^α at points 1 and 2 involves a series of terms. In Fig. 1 we have listed the bare loop and tree graphs with scalar condensates up to dimension 6. If the calculation is performed in an arbitrary gauge, additional graphs like those in Fig. 2 have to be included. In the case of the Landau gauge, these contributions vanish due to the orthogonality of the projection operators,

$$P_T P_i = P_T P_2 = 0 . \quad (38)$$

The graphs of Fig. 1 yield, for the Fourier-transformed two-point functions, the simple result (cf. Eqs. (28), (30), (31)):

$$\begin{aligned} & i \int d^4x e^{iqx} \langle T(J_\beta^\mu(x, \theta, \bar{\theta}) J_\delta^\nu(0, 0, 0)) \rangle \\ &= \frac{1}{2} P_T \delta^4(\theta) q^2 \Pi_{\beta\delta}^{\mu\nu}(q^2) - \frac{1}{2} (P_i - P_2) \delta^4(\theta) \delta_\delta^\mu \delta_\beta^\nu (\langle \chi^{*\mu} \chi_\delta \rangle - \langle \chi^{*\nu} \chi_\beta \rangle) \\ & \quad - \frac{1}{2} \delta^4(\theta) \delta_\delta^\mu \delta_\beta^\nu (\langle \chi^{*\mu} \chi_\delta \rangle + \langle \chi^{*\nu} \chi_\beta \rangle) , \end{aligned} \quad (39a)$$

with

$$\begin{aligned}
 \Pi_{\beta\delta}^{\alpha\gamma}(q^2) &= \delta_{\delta}^{\alpha} \delta_{\beta}^{\gamma} \left[-\frac{1}{8\pi^2} \ln \frac{q^2}{\Lambda^2} + \frac{1}{q^2} (\langle \chi^* \chi_{\alpha} \rangle + \langle \chi^* \chi_{\beta} \rangle) \right] \\
 &\quad - 2 \frac{2\pi\alpha}{(q^2)^2} \langle (\chi^* \tau^A \chi_{\beta}) (\chi^* \tau^A \chi_{\delta}) \rangle \\
 &\quad + 2 \frac{(2\pi\alpha)^2}{(q^2)^3} \langle (\chi^* \tau^A \chi_{\beta}) (\chi^* \tau^A \chi_{\delta}) (\chi^* \chi) \rangle + \dots, \\
 \alpha &= \frac{g^2}{4v}, \quad \chi^* \chi \equiv \sum_{\epsilon} \chi^* \chi_{\epsilon}.
 \end{aligned} \tag{39b}$$

For a conserved current J_{β}^{α} one naively expects

$$\partial^2 \langle T(J_{\beta}^{\alpha}(x, \theta, \bar{\theta}) J_{\delta}^{\gamma}(0, 0, 0)) \rangle = 0 \tag{40}$$

Equation (40) is violated due to the terms in (39a) which are not proportional to P_T . These terms correspond to the contribution of Goldstone chiral multiplets to the two-point functions of the currents J_{β}^{α} and are analogous to Schwinger terms which occur in two-point functions of currents involving derivatives (cf. Ref. 10)). We expect that the transversality of the expression (39a) can be restored by using a properly defined T^* product.

The tensor $\Pi_{\beta\delta}^{\alpha\gamma}(q^2)$ contains, with respect to $SU(2) \times U(4)$, invariants which transform as (3,1), (1,1), (1,15), (1,1)' and (2,4):

$$\Pi_{ij}^{kl} = (\delta_{\ell}^i \delta_j^k - \frac{1}{2} \delta_j^i \delta_{\ell}^k) \Pi^{(3,1)} + \frac{1}{2} \delta_j^i \delta_{\ell}^k \Pi^{(1,1)}, \tag{41a}$$

$$\Pi_{bd}^{ac} = (\delta_d^a \delta_b^c - \frac{1}{4} \delta_b^a \delta_d^c) \Pi^{(1,15)} + \frac{1}{4} \delta_b^a \delta_d^c \Pi^{(1,1)'}, \tag{41b}$$

$$\Pi_{aj}^{ib} = \delta_j^i \delta_a^b \Pi^{(2,4)}. \tag{41c}$$

From Eqs. (29), (32), (39) and (41) one obtains:

$$\begin{aligned} \overline{\Pi}^{(3,1)}(q^2) = & -\frac{1}{8\pi^2} \ln \frac{q^2}{\Lambda^2} + \frac{1}{q^2} \langle \chi^\dagger \chi \rangle - \frac{2\overline{u}d}{(q^2)^2} \langle (\chi^\dagger \overline{\tau} \chi)^2 \rangle \\ & + \frac{(2\overline{u}d)^2}{(q^2)^3} \langle (\chi^\dagger \overline{\tau} \chi)^2 (\chi^\dagger \chi) \rangle + \dots, \end{aligned} \quad (42a)$$

$$\overline{\Pi}^{(1,1)}(q^2) = -\frac{1}{8\pi^2} \ln \frac{q^2}{\Lambda^2} + \frac{1}{q^2} \langle \chi^\dagger \chi \rangle, \quad (42b)$$

$$\overline{\Pi}^{(2,4)}(q^2) = -\frac{1}{8\pi^2} \ln \frac{q^2}{\Lambda^2} + \frac{1}{2q^2} \langle \chi^\dagger \chi \rangle, \quad (42c)$$

$$\begin{aligned} \overline{\Pi}^{(1,15)}(q^2) &= \overline{\Pi}^{(1,1)'}(q^2) \\ &= -\frac{1}{8\pi^2} \ln \frac{q^2}{\Lambda^2}. \end{aligned} \quad (42d)$$

In channels without Goldstone bosons it is reasonable to argue¹⁰⁾ that $P_T \delta(1,2)$ is the right covariant to be taken out of the two-point function like in the case of the free vector propagator. Then $\Pi^{(3,1)}$, $\Pi^{(1,15)}$ and $\Pi^{(1,1)'}$ have to be changed to $\hat{\Pi} = q^2 \Pi$. This matters particularly in the (3,1) channel where the $\langle \bar{\chi} \chi \rangle$ term does not contribute to $\text{Im } \hat{\Pi}^{(3,1)}$. In the following analysis, however, this point is not important.

For the invariant functions $\Pi(q^2)$ we can now consider the familiar Borel sum rules⁵⁾:

$$\begin{aligned} \frac{1}{\Pi} \int_0^\infty ds \text{Im } \overline{\Pi}_P(s) e^{-\frac{s}{M^2}} \\ = M^2 \left[h_0 + \frac{h_1}{M^2} + \frac{1}{2!} \frac{h_2}{M^4} + \frac{1}{3!} \frac{h_3}{M^6} + \dots \right], \end{aligned} \quad (43a)$$

with

$$q^2 \left(-\frac{d}{dq^2} \right) \overline{\Pi}(q^2) = h_0 + \frac{h_1}{q^2} + \frac{h_2}{(q^2)^2} + \frac{h_3}{(q^2)^3} + \dots, \quad (43b)$$

where M is the Borel variable and $\text{Im } \Pi_P(s)$ the physical spectral function.

Given a certain Ansatz for the physical spectrum, the sum rules (43) can be used to constrain the resonance parameters. In analogy to QCD one would expect

that the dominant contributions come from the lowest-lying vector multiplets which are created from the vacuum by the currents J_β^α , and - in the $SU(2) \times U(4)$ phase - the Goldstone multiplets. In addition there is, contrary to QCD, a massless two-particle continuum⁷⁾ to which, in the $U(6)$ phase, all 15 chiral multiplets contribute. In the broken phase only the 9 Goldstone multiplets ϕ_{ia} and ϕ remain massless; the $SU(4)$ sextet ϕ_{ab} acquires a mass of order Λ through the instanton process shown in Fig. 3 which can be calculated along the lines of Ref. 19). The contribution of the massless physical continuum at $s = 0$ is easily obtained by expressing the $U(6)$ currents in terms of the composite chiral multiplets,

$$J_\beta^\alpha = \bar{\Phi}^{\alpha\gamma} \Phi_{\beta\gamma} + \dots \quad (44)$$

In the case of unbroken $U(6)$ symmetry we are thus led to the following Ansatz for the spectral functions of the two irreducible $SU(6)$ representations:

$$\frac{1}{11} \text{Im} \bar{\Pi}^{(35)}(s) = \frac{1}{4\pi^2} \Theta(t_u - s) + \frac{2M_u^2}{g_u^2} \delta(s - M_u^2) \quad , \quad (45a)$$

$$\frac{1}{11} \text{Im} \bar{\Pi}^{(1)}(s) = \frac{6}{8\pi^2} \Theta(t_{u'} - s) + \frac{2M_{u'}^2}{g_{u'}^2} \delta(s - M_{u'}^2) \quad , \quad (45b)$$

where t_u and $t_{u'}$ are the scales up to which the massless chiral multiplets couple pointlike to the currents J_β^α . The usual heavy continuum has been neglected because in the following we will only be interested in the resonance region and not in the asymptotic behaviour of the spectral functions for large values of s . In the $SU(2) \times U(4)$ phase the sextet ϕ_{ab} does not contribute to the massless continuum and instead of Eqs. (45) one arrives at the Ansatz:

$$\frac{1}{11} \text{Im} \bar{\Pi}^{(3,1)}(s) = \frac{1}{4\pi^2} \Theta(t - s) + \frac{2M_w^2}{g_w^2} \delta(s - M_w^2) \quad , \quad (46a)$$

$$\frac{1}{11} \text{Im} \bar{\Pi}^{(1,1)}(s) = \frac{3}{8\pi^2} \Theta(t - s) + 2f_1^2 \delta(s) + \frac{2M_{w'}^2}{g_{w'}^2} \delta(s - M_{w'}^2) \quad , \quad (46b)$$

$$\frac{1}{11} I_{\mu} \Pi^{(2,4)}(s) = \frac{1}{16\pi^2} \Theta(t-s) + f_2^2 \delta(s) + \frac{2M_Y^2}{g_Y^2} \delta(s-M_Y^2) , \quad (46c)$$

$$\frac{1}{11} I_{\mu} \Pi^{(1,15)}(s) = \frac{1}{8\pi^2} \Theta(t-s) + \frac{2M_u^2}{g_u^2} \delta(s-M_u^2) , \quad (46d)$$

$$\frac{1}{11} I_{\mu} \Pi^{(1,1)'}(s) = \frac{1}{8\pi^2} \Theta(t-s) + \frac{2M_{u'}^2}{g_{u'}^2} \delta(s-M_{u'}^2) . \quad (46e)$$

$(1,1)$ and $(1,1)'$ denote the $U(1)$ factors of the $U(2)$ and $U(4)$ subgroups of $U(6)$. The terms proportional to f_1^2 and f_2^2 are due to the direct coupling of the massless Goldstone multiplets to the currents of the spontaneously broken symmetries.

Let us first consider the case of unbroken $U(6)$ symmetry. From Eqs. (42), (43) and (45) one obtains

$$\frac{1}{4\pi^2} M^2 (1 - e^{-\frac{t_u}{M^2}}) + \frac{2M_u^2}{g_u^2} e^{-\frac{M_u^2}{M^2}} = \frac{1}{8\pi^2} M^2 , \quad (47a)$$

$$\frac{5}{8\pi^2} M^2 (1 - e^{-\frac{t_{u'}}{M^2}}) + \frac{2M_{u'}^2}{g_{u'}^2} e^{-\frac{M_{u'}^2}{M^2}} = \frac{1}{8\pi^2} M^2 . \quad (47b)$$

Equations (47) are very different from QCD sum rules for hadronic resonances⁵⁾. The main difference lies in the absence of power corrections to the perturbative continuum. This is the case also in higher orders in α because, as we saw in the previous section, all condensates vanish in the unbroken $U(6)$ phase. Therefore there is no scale which could determine the resonance masses. Furthermore it is disturbing that the massless bound state continuum is larger than the contribution from the constituents. Thus the Ansatz, which yields Eqs. (47), seems very unnatural, and one is led to the conclusion that the phase of unbroken $U(6)$ symmetry is unlikely to exist at all. This result is in agreement with "complementarity" as we have seen in Section 2.

In the $SU(2) \times U(4)$ case, Eqs. (42), (43) and (46) imply the sum rules:

$$(3,1): \frac{1}{4\pi^2} M^2 (1 - e^{-\frac{t_W}{M^2}}) + \frac{2M_W^2}{g_W^2} e^{-\frac{M_W^2}{M^2}} \quad (48a)$$

$$= \frac{1}{8\pi^2} M^2 + \langle \chi^* \chi \rangle - \frac{2\pi d}{M^2} \langle (\chi^* \bar{\chi})^2 \rangle + \frac{1}{2!} \frac{(2\pi d)^2}{M^4} \langle (\chi^* \bar{\chi})^2 (\chi^* \chi) \rangle + \dots$$

$$(1,1): \frac{3}{8\pi^2} M^2 (1 - e^{-\frac{t_{W'}}{M^2}}) + 2f_1^2 + \frac{2M_{W'}^2}{g_{W'}^2} e^{-\frac{M_{W'}^2}{M^2}} = \frac{1}{8\pi^2} M^2 + \langle \chi^* \chi \rangle, \quad (48b)$$

$$(2,4): \frac{1}{16\pi^2} M^2 (1 - e^{-\frac{t_Y}{M^2}}) + f_2^2 + \frac{2M_Y^2}{g_Y^2} e^{-\frac{M_Y^2}{M^2}} = \frac{1}{8\pi^2} M^2 + \frac{1}{2} \langle \chi^* \chi \rangle, \quad (48c)$$

$$(1,15): \frac{1}{8\pi^2} M^2 (1 - e^{-\frac{t_U}{M^2}}) + \frac{2M_U^2}{g_U^2} e^{-\frac{M_U^2}{M^2}} = \frac{1}{8\pi^2} M^2, \quad (48d)$$

$$(1,1)': \frac{1}{8\pi^2} M^2 (1 - e^{-\frac{t_{U'}}{M^2}}) + \frac{2M_{U'}^2}{g_{U'}^2} e^{-\frac{M_{U'}^2}{M^2}} = \frac{1}{8\pi^2} M^2. \quad (48e)$$

Contrary to the $U(6)$ case, Eqs. (48) contain contributions from scalar condensates in all channels in which the massless continua due to constituents and bound states do not match. These power corrections appear with large coefficients because the scalar condensates contribute through tree graphs. The sum rule in the $(3,1)$ channel is essentially identical with the one obtained previously¹⁰⁾ for the $SU(2)$ current in the Abbott-Farhi model. In the channels $(1,15)$ and $(1,1)'$, small power corrections will be generated at higher orders in α . The leading contribution is of order $\alpha^2/\pi^2 \langle \chi^* \chi \rangle$ (cf. Fig. 4).

The consequences which can be drawn from the sum rules depend on the size of the condensate $\langle \chi^* \chi \rangle$, i.e., on the value of $\alpha (\langle \chi^* \chi \rangle)$. In a confining theory with a single mass scale, one expects $\alpha (\langle \chi^* \chi \rangle) \sim 1$. The perturbative expansion in α , which leads to Eqs. (48), is then consistent for values of M in the range $8\pi^2 > M^2 / \langle \chi^* \chi \rangle > 1$. This means that contrary to QCD there exists a region in M^2 where power "corrections" dominate the perturbative continuum and perturbation theory in α is still reliable.

The large power corrections in the $(3,1)$ channel are clearly related to the W -mass. Assuming, for simplicity, the factorization¹⁰⁾

$$\langle (\chi^* \bar{\chi})^2 (\chi^* \chi)^n \rangle = \langle \chi^* \chi \rangle^{n+2}, \quad (49)$$

they can be summed up yielding the standard model mass formula¹⁰⁾

$$M_W^2 = 2\bar{u} \alpha_W \langle \chi^* \chi \rangle, \quad (50)$$

as well as $\alpha(M^2) = \alpha_W(M^2)$. Due to the absence of higher tree graph power corrections in all other channels, it seems natural to assume $M_{W'} \approx M_Y \approx M_{U'}$, and $g_{W'} \approx g_Y \approx g_{U'} \approx g_U$, which implies

$$f_1^2 \approx f_2^2 \approx \frac{1}{2} \langle \chi^* \chi \rangle, \quad (51)$$

and

$$\left| \frac{f_2^2 - f_1^2}{f_2^2} \right| \approx \frac{M^2}{8\bar{u} \langle \chi^* \chi \rangle} \ll 1. \quad (52)$$

The parameters of the vector bosons $W', \dots, U'$ are not directly related to the scalar condensates. If these vector bosons exist, they must interact strongly. With $M = M_{W'}$, one obtains, as for the ρ -meson⁵⁾

$$d_{W'} = \frac{g_{W'}^2}{4\pi} \approx \frac{4\pi}{e}, \quad (53)$$

where e is Euler's constant. According to the effective Lagrangian (1), the ratio $(f_2^2 - f_1^2)/f_2^2$ is equal to the ratio of isovector and isoscalar coupling strengths, i.e.,

$$\left| \frac{f_2^2 - f_1^2}{f_2^2} \right| \approx \frac{d_{W'}}{M_{W'}^2} \frac{M_W^2}{d_W}. \quad (54)$$

From Eqs. (52) - (54) we then obtain that the vector bosons $W', \dots, U'$ are substantially heavier than the W -bosons,

$$M_{W'}^2 \gg \frac{4\pi}{e} \frac{M_W^2}{d_W} > \frac{4\pi}{e} M_W^2, \quad (55)$$

where we have used $\alpha_W(M^2) = \alpha(M^2) < 1$.

Although the sum rules (48) as well as the low energy effective Lagrangian (1) are compatible with the existence of strongly interacting heavy vector bosons, the simplest Ansatz satisfying Eqs. (48) requires only the states present

in the complementary picture, i.e., the Goldstone multiplets ϕ_{1a} and ϕ , and the SU(2) triplet of W-bosons. At $M^2 = M_W^2 = 2\pi\alpha \langle \chi^* \chi \rangle$, the differences between the perturbative and the bound state continuum contributions in the channels (3,1), (1,1) and (2,4) can be accounted for by radiative corrections.

5. CONCLUSIONS

In the previous sections we have studied scalar condensates and their influence on the physical spectrum in a supersymmetric model with SU(2) gauge interaction and global U(6) invariance. The model corresponds to SQCD with two colours and three flavours in the chiral limit. The mass dependence of bilinear condensates, as it is suggested by the effective Lagrangian approach and the SUSY Dashen formulae, implies the following vacuum structure: either all condensates vanish, in which case the U(6) symmetry remains unbroken, or U(6) is spontaneously broken to SU(2) $\times$ U(4).

Given the condensates of a theory, one can study their effect on the physical spectrum by means of SVZ sum rules. We have computed the dominant contributions of the scalar condensates to the two-point functions of the U(6) currents. The result indicates that there is no phase with unbroken U(6) symmetry. The sum rules obtained in the SU(2) $\times$ U(4) case are qualitatively different from QCD sum rules: the power corrections are large because the scalar condensates contribute through tree graphs, and can even dominate the perturbative continuum; furthermore there is a contribution to the continuum from massless bound states. The obtained sum rules are compatible with the existence of a SU(2) triplet of W-bosons and additional 33 vector bosons $W', \dots, U'$, which are heavier and more strongly interacting than the W-bosons. The simplest and most natural Ansatz satisfying the sum rules contains, however, only the states of the complementary picture, i.e., the Goldstone multiplets and the W-vector multiplets. We thus arrive at the conclusion that, contrary to naive expectations, the physical spectra in the confining phase and the Higgs phase may be identical.

In a recent paper²⁴⁾, Kugo, Uehara and Yanagida have obtained the supersymmetric σ -model for the coset space U(6)/SU(2) $\times$ U(4) in closed form. They find that the Lagrangian possesses a local SU(2) invariance and suggest the existence of a corresponding SU(2) triplet of massive dynamical gauge bosons. This suggestion, in particular the special role of the SU(2) factor of the unbroken SU(2) $\times$ U(4) symmetry, is in agreement with our results.

The motivation of our investigation was to understand whether it is possible to consider the intermediate vector bosons of the weak interactions as composite vector bosons in a preon theory of quarks and leptons. As we discussed in the introduction, this requires that the mass of the W-bosons is small compared to the typical vector boson masses of the theory and that their coupling to the SU(2) current is weak, $\alpha_W \sim 1/20$. It is conceivable that the first condition can be satisfied by means of appropriate scalar condensates. To realize the second condition appears much more difficult. Our analysis suggests $\alpha_W(M_W^2) \approx \alpha(M_W^2) \approx 1$ where $g = \sqrt{4\pi\alpha}$ is the $SU(2)_{HC}$ gauge coupling constant. In order to avoid this problem one has to invoke either rather unconventional factorization properties of higher dimensional condensates¹⁰⁾ or the scalar condensates have to be very large, i.e., $\alpha(\langle\chi^*\chi\rangle) \ll 1$. The case of large condensates, however, is by definition the Higgs phase of the theory where our analysis is merely a gauge invariant reformulation of the standard model. We have to conclude that, at least within supersymmetric models, it appears difficult to obtain the W-bosons of the weak interactions as composite vector bosons.

Our analysis led to the unexpected result that the spectra of physical states in the Higgs phase and the confining phase may be identical. This suggests that the intuition derived from QCD may be very misleading in the case of SQCD and emphasizes the importance of the complementary picture.

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FIGURE CAPTIONS

- Fig. 1 Contributions to the two-point functions of the current J_{β}^{α} in Landau gauge.

- Fig. 2 Additional contributions to the two-point functions of the currents J_{β}^{α} in an arbitrary gauge.

- Fig. 3 Instanton process which generates a mass for the SU(4) sextet ϕ_{ab} in the case of unbroken SU(2) $\times$ U(4) symmetry.

- Fig. 4 Contribution to the dominant power correction in the (1,15) and (1,1)' channels.

Fig. 1

Fig. 2

Fig. 3

Fig. 4