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ON THE EXCITED QUARKS AND LEPTONS

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A B S T R A C T

Using a duality-like finite energy sum rule, we discuss the assumption of having excited fermions at the W-scale in a supersymmetric (SUSY) and non-supersymmetric hypercolour theory where quarks and leptons are bound states of fermion and scalar preon constituents. We conclude that a SUSY-like composite model cannot have excited fermions having a mass smaller than 0.5 TeV. A non-SUSY composite model having composite fermions but elementary W-bosons can produce an excited fermion mass of the order of M_W provided that the scalar vacuum condensate is of the order of the $(\text{TeV})^2$ scale of compositeness.

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1. INTRODUCTION

We have studied in Refs. 1) and 2) constraints among the spin 1 and 0 bound states, the continuum threshold and the lowest dimension vacuum condensates of a class of composite models of weak interactions which are described at the TeV scale by an hypercolour vector-like theory. We found that in the non-supersymmetric case, the dynamics of the theory is dominated by the large value of the hypergluon vacuum condensate $\alpha_H \langle F^2 \rangle$ which characterizes the difference of the hypercolour theory with the usual QCD theory, whilst the vacuum condensate of scalar fields cannot be as large as the W-boson mass squared, unless there exists a light excitation of the W-boson. On the other hand, it has been shown²⁾ that the spin zero partners of the W-bosons having a mass of the order of 50 GeV^{*}) can be understood with a relatively small explicit breaking of the preon chiral symmetry provided that these spin zero particles are not of the η' -type.

We extend the analysis in Refs 1) and 2) to the case of the spin $\frac{1}{2}$ bound states built from fermion and scalar elementary fields.

The possibility for having composite quarks and leptons is very appealing as in this way, one may dynamically understand the relatively small values of the quark and lepton masses⁴⁾. According to the anomaly matching condition⁵⁾, the lightness of these masses is due to an unbroken chiral symmetry at the constituent preon level. In the quasi-Goldstone scheme⁶⁾, quarks and leptons are partners of some Goldstone bosons associated to the breakdown of a global symmetry group G_f into H_f , so that they belong to the adjoint of G_f/H_f . Unwanted massless Goldstone bosons may get masses through a breakdown of supersymmetry (SUSY). In the haplon scheme⁷⁾, the fermion masses can be generated by a radiative correction mechanism^{4a), 8)}. In both schemes, the masses of the fermions are light compared to the compositeness scale of fermions which is known to be of the order of TeV from the analysis of the deviation of the QED-like process due to the quark form factor or (and) to an additional four-fermion type interaction⁹⁾.

) This value of the mass is the one suggested by the UA1,2 data for the decay $Z \rightarrow e^+ e^- \gamma$ ³⁾, if it occurs via a virtual e^ .

One would naturally expect from a duality-like sum rule approach^{10),1),2)*)} that the level spacing between the lowest ground state, i.e. the quark and leptons and the barycentre of the first two or three excitation masses is of the same order as the compositeness scale. However, due to the Goldstone nature of fermions, one might expect the particular situation where the duality argument may not work rigorously as the dominance of the spectral function by the lowest ground state fermion will be lost at the TeV scale where one expects that the QHD asymptotic description of the weak interactions starts to make sense, requiring the first excited state to have a mass much lower than the compositeness scale so that it will give an appreciable contribution to the spectral function.

From a phenomenological point of view, the possibility for having excited fermions of mass around 70-160 GeV can give one explanation¹²⁾ (among others¹³⁾) of various unexpected events in the SppS experiments (large rate of $Z \rightarrow e^+ e^- \gamma$, for the explanation of the $Z \rightarrow e^+ e^- \gamma$ events requires a definite chirality for the e^* , otherwise it will induce an anomalously large contribution to the $(g-2)^{14)}$ which is unwanted from the present agreement between the theoretical calculations¹⁵⁾ and the experimental number¹⁶⁾). In this paper, we shall be interested in the question of whether the possibility to have an excited electron at the W-scale can be understood dynamically. Our approach will be based on QCD-like spectral function sum rules which have been used earlier in Ref. 10c) for the Laplace transform sum rule¹⁷⁾ and later on ^{1),2)**)} within the finite energy sum rule (FESR) context²⁰⁾. The method allows to relate the bound states and the continuum parameters to the dynamical parameters of the models which are assumed to be described by the vacuum condensates. We shall work for definiteness in

*) The unorthodox case which is inconsistent with the duality approach is discussed in Ref.11). In the duality framework, the lowest resonance absorbs into its effective contribution the one coming from the continuum having a threshold below the resonance mass. This is analogous to the case of the ρ meson and of the $\pi^+-\pi^-$ continuum threshold which starts below the ρ -meson mass.

***) A prior work using FESR within a slightly different context is in Ref. 18). The role of the scalar condensate in a theory with scalar fields has been noticed first by the authors in Ref. 18) and later on by many others^{1),19)}.

the class of models where the fermion is built with a scalar field $\phi^{*})$ and a Weyl Spinor ϕ_L in such a way that the fermion is chiral, because of our ignorance of the exact nature of fermions.

2. NON-SUPERSYMMETRIC MODEL

We work for definiteness with a QHD Lagrangian based on the hypercolour group $SU(N)_H$ which contains a fermion and a scalar fields. The Lagrangian density is:

$$\begin{aligned} \mathcal{L}(x) = & -\frac{1}{4} F_{\mu\nu}^a(x) F_a^{\mu\nu}(x) + \int \frac{d^4p}{(2\pi)^4} (\bar{\mathcal{D}}_\mu^* \phi_f^+) (\mathcal{D}^\mu \phi_f) - V(\phi^+ \phi) \\ & + \frac{i}{2} \int \frac{d^4p}{(2\pi)^4} \bar{\Psi}_f(x) \delta^{\mu\leftrightarrow} \mathcal{D}_\mu \Psi_f(x) - \int \frac{d^4p}{(2\pi)^4} m_f \bar{\Psi}_f \Psi_f \end{aligned} \quad (1)$$

which is invariant under the $SU(N)_H$ transformations; a and f are the hypercolour and preon indices; $\phi(x)$ and $\phi(x)$ are respectively the spinor and complex spin zero fields in the fundamental representations; the potential V is the usual scalar potential; $\mathcal{D}^\mu$ is the covariant derivative; $F_{\mu\nu}$ is the hypergluon strength tensor. We shall be interested in the analysis of the two-point functions:

$$\begin{aligned} S_f(q) &= i \int d^4x e^{iqx} \langle 0 | \mathbb{T} J_f(x) \bar{J}_f(0) | 0 \rangle, \\ \pi^{\mu\nu}(q) &= i \int d^4x e^{iqx} \langle 0 | \mathbb{T} V^\mu(x) (V^\nu(0))^+ | 0 \rangle, \end{aligned} \quad (2)$$

associated respectively to the spinorial and vector currents:

$$J_f(x) \equiv \phi_f^+ \frac{1}{2} (1 - \sigma_5) \Psi \quad , \quad (3)$$

$$V^\mu(x) \equiv \bar{\Psi} \gamma^\mu \Psi + i \phi_f^+ \overleftrightarrow{\mathcal{D}}^\mu \phi_f, \quad (4)$$

where the expression in Eq. (4) is the unique combination derived from the equation of motion $\delta L / \delta A_\mu^a = 0$ (A_μ^a is the Hypergluon field) which respects the transversality of $\pi^{\mu\nu}(q)$ in the theory, having both spinor and scalar elementary fields. In the case studied here, the current in Eq. (3) will correspond to the fermion bound state while the one in Eq. (4) will correspond to the W-bosons if they are composite objects. We study the asymptotic behaviour of these two-

*) The scalar field ϕ can be the "dynamical scalar" [according to the terminology of Ref. 21)] formed by two elementary spin $\frac{1}{2}$ preons which can give a "natural" explanation of the lightness of quarks and leptons.

point functions within the Wilson's operator product expansion (OPE) of SVZ^{17a)} which incorporates in a systematic way the non-trivial contributions of the vacuum condensates of the QHD theory. Their QHD expressions come from the evaluation of the QHD diagrams in Fig. 1 and 2 and read respectively^{*)}:

$$S_q(q) = \frac{1}{2} (1 - \delta_5) \hat{q} \left\{ -\frac{N}{32\pi^2} \log \frac{Q^2}{\nu^2} + \frac{\langle \phi^\dagger \phi \rangle}{Q^2} - \frac{\alpha_H \langle F^2 \rangle}{64\pi Q^4} + \dots \right\}, \quad (5)$$

$$\begin{aligned} \Pi^{\mu\nu}(q) = & - (q^\mu q^\nu - q^2 g^{\mu\nu}) \left\{ -\frac{5N}{48\pi^2} \log \frac{Q^2}{\nu^2} + 2 \frac{\langle \phi^\dagger \phi \rangle}{Q^2} \right. \\ & \left. - 8 \left(\frac{N^2 - 1}{N^2} \right) \pi \alpha_H \frac{\langle \phi^\dagger \phi \rangle^2}{Q^4} + \frac{\alpha_H \langle F^2 \rangle}{64\pi Q^4} + \dots \right\} \quad (6) \\ & (Q^2 = -q^2 > 0) \end{aligned}$$

The contributions of the scalar condensate are straightforward to obtain and agree with previous results^{1), 18), 19a), c)}. The $\alpha_H \langle F^2 \rangle$ hypergluon condensate in Eq. (6) agrees with the one in Ref. 18). We relate the QHD expression to the bound state parameters using the FESR technique à la QCD²⁰⁾ which appears to be more convenient for the study of the composite model dynamics^{1), 2), 18)} than the Laplace transform sum rule¹⁷⁾ because of the many uncontrollable unknowns for the dynamics of the model. Parametrizing the lepton or quark propagators as:

$$S_q(q) = \frac{1}{2} (1 - \gamma_5) \hat{q} \left\{ \sum_{i=1}^N \frac{\lambda_i^2}{q^2 + m_i^2} + \text{"QHD continuum"} \right\}, \quad (7)$$

where λ_i is the lepton's wave function residue, the FESR analysis gives the set of constraints:

$$\sum_{i=e, e^*} \lambda_i^2 = \frac{N}{32\pi^2} t_c \approx \langle \phi^\dagger \phi \rangle, \quad (8a)$$

$$\sum_{i=e, e^*} \lambda_i^2 m_i^2 = \frac{N}{64\pi^2} t_c^2 \approx \frac{\alpha_H \langle F^2 \rangle}{64\pi}, \quad (8b)$$

where t_c is the QHD continuum threshold which is expected to be of the order of the TeV scale of compositeness. We use the positivity of λ_e^2 in Eq. (8a) and we retain the leading term in m_e^2 in Eq. (8b). Then one gets for $SU(2)_H$ the constraint:

*) We give in the Table the different contributions of each Feynmann diagrams in the case of SUSY theory. The one for non-SUSY can be obtained from the rules in the table caption.

$$m_{e^*}^2 \gtrsim \frac{t_c}{2} \left(\frac{1 + \frac{\pi \alpha_H \langle F^2 \rangle}{2 t_c^2}}{1 + 16 \pi^2 \langle \phi^+ \phi \rangle / t_c} \right). \quad (9)$$

In the model having elementary W-bosons, the results in Eq. (6) cannot provide any useful information. For this case, the assumption of having a e^* -mass at the W scale requires a large value of the scalar vacuum condensate:

$$\langle \phi^+ \phi \rangle \gtrsim \left(\frac{t_c^2 + \frac{\pi}{2} \alpha_H \langle F^2 \rangle}{32 \pi^2 m_{e^*}^2} \right) \simeq (0.8 \text{ TeV})^2 \quad (10)$$

where the numerical value corresponds to $t_c \simeq 1 \text{ TeV}^2 \simeq \alpha_H \langle F^2 \rangle$ and $m_{e^*} \simeq 70 \text{ GeV}$.

In the case where the model has composite W-bosons, Eq. (6) gives the constraint for $SU(2)_H$:

$$F_W^2 - \frac{5 t_c}{24 \pi^2} \simeq 2 \langle \phi^+ \phi \rangle \quad (11a)$$

$$F_W^2 M_W^2 - \frac{5 t_c^2}{48 \pi^2} \simeq 6 \pi \alpha_H \langle \phi^+ \phi \rangle^2 - \frac{\alpha_H \langle F^2 \rangle}{6 \pi}, \quad (11b)$$

which is very similar to the one discussed in Ref. 1) for the case of the spin 0 constituents of the W-bosons. We know from such an analysis that the scalar condensate should be smaller than M_W^2 from the positivity of t_c while the value of the hypergluon vacuum condensate might be of the order of the TeV^4 . One can see, in this case, that the e^* -mass is larger than $\sqrt{t_c}/2$, i.e., there is no room to have a e^* -mass at the W-scale for this class of models where both the fermions and the W-bosons are composite objects.

3. SUPERSYMMETRIC MODEL

The supersymmetric version of the Lagrangian in Eq. (1) is equivalent to the Lagrangian of supersymmetric QCD which can be obtained in existing literatures²²⁾. For our purpose, the equivalent of Eq. (3) can be obtained from the θ -component of the superpotential $m \phi \tilde{\phi}$ where respectively ϕ and $\tilde{\phi}$ are chiral and anti-chiral matter superfields:

$$\begin{aligned} \varphi &\equiv \phi + \sqrt{2} \theta \psi + \dots \\ \tilde{\varphi} &\equiv \tilde{\phi} + \sqrt{2} \theta \tilde{\psi} + \dots \end{aligned} \quad (12)$$

where θ is an anti-commuting variable. Then, in a two-component notation:*)

$$J_q(x) \equiv \frac{1}{\sqrt{2}} (\phi \tilde{\psi} + \psi \tilde{\phi}) . \quad (13)$$

In the same way, one can derive the vector current from the interaction term $\frac{1}{2} (\lambda_a/2) V_a \phi \tilde{\phi}$ where V_a is the vector superfield²²⁾. Then, in a two-component notation:

$$J_\mu(x) = \tilde{\psi} \sigma_\mu \tilde{\bar{\psi}} - \psi \sigma_\mu \bar{\psi} + i (\phi^+ \tilde{D}^\mu \phi + \tilde{\phi}^+ \tilde{D}^\mu \tilde{\phi}) . \quad (14)$$

The new feature of supersymmetry is the existence of various condensates induced by the Yukawa coupling of $g\lambda\phi\phi$ where λ is the spin $\frac{1}{2}$ partner of the hypergluon field. For instance we show in Fig. 3 the new additional contribution to $S_F(q)$. Then, the expressions of the associated two-point function in the limit of the zero mass superfield are (see also the Table):

$$S_q(q) = \left\{ -\frac{N}{32\pi^2} \log \frac{q^2}{\Lambda^2} + \frac{1}{2} \langle \phi^+ \phi + \tilde{\phi}^+ \tilde{\phi} \rangle \frac{1}{q^2} - \frac{3}{8} \sqrt{2} g_H \langle \phi^+ \lambda \psi + \tilde{\phi}^+ \lambda \tilde{\psi} \rangle \frac{1}{q^4} - \frac{\alpha_H \langle F^2 \rangle}{64\pi} \frac{1}{q^4} - \frac{3}{4} \left(\frac{N^2-1}{N^2} \right) \pi \alpha_H \langle \phi^+ \phi + \tilde{\phi}^+ \tilde{\phi} \rangle^2 \frac{1}{q^4} - \frac{1}{4} \left(\frac{N^2-1}{N^2} \right) \pi \alpha_H \langle \phi \tilde{\phi} + \phi^+ \tilde{\phi}^+ \rangle \frac{1}{q^4} \right\} \frac{1}{2} (1-\gamma_5) \not{q} , \quad (15)$$

$$\pi^{\mu\nu}(q) = - (q^2 g^{\mu\nu} - q^\mu q^\nu) \left\{ -\frac{N}{8\pi^2} \log \frac{q^2}{\Lambda^2} + 2 \langle \phi^+ \phi + \tilde{\phi}^+ \tilde{\phi} \rangle \frac{1}{q^2} - 2\sqrt{2} g_H \langle \phi^+ \lambda \psi + \tilde{\phi}^+ \lambda \tilde{\psi} \rangle \frac{1}{q^4} + \frac{\alpha_H \langle F^2 \rangle}{4\pi} \frac{1}{q^4} - 2 \left(\frac{N^2-1}{N^2} \right) \pi \alpha_H \langle \phi^+ \phi + \tilde{\phi}^+ \tilde{\phi} \rangle^2 \frac{1}{q^4} - 6 \left(\frac{N^2-1}{N^2} \right) \pi \alpha_H \langle \phi \tilde{\phi} + \phi^+ \tilde{\phi}^+ \rangle \frac{1}{q^4} \right\} , \quad (16a)$$

with:

$$\langle \phi^+ \lambda \psi + \tilde{\phi}^+ \lambda \tilde{\psi} \rangle \equiv -i \langle 0 | \phi^+ \epsilon_a \lambda^a \psi + \tilde{\phi}^+ \epsilon_a \lambda^a \tilde{\psi} | 0 \rangle . \quad (16b)$$

*) We adopt the Weyl notation for the spinors:

$$\psi_D \equiv \begin{pmatrix} \psi_\alpha \\ \tilde{\bar{\psi}}^{\dot{\alpha}} \end{pmatrix} ; \quad \bar{\psi}_D \equiv (\tilde{\psi}^\alpha, \bar{\psi}_{\dot{\alpha}})$$

The fact that supersymmetry is unbroken by the condensates implies the following relation:

$$\sqrt{2}g_H \langle \phi^+ \lambda \psi + \tilde{\phi} \lambda \tilde{\psi} \rangle = \left(\frac{N^2-1}{N^2} \right) \pi \alpha_H \left\{ \langle \phi^+ \phi + \tilde{\phi}^+ \tilde{\phi} \rangle^2 - \langle \phi \tilde{\phi} + \phi^+ \tilde{\phi}^+ \rangle^2 \right\}, \quad (17)$$

where we have used a factorization hypothesis of the four-scalar condensate operator. Using the fact that:

$$\langle \phi^+ \phi \rangle \simeq \langle \tilde{\phi}^+ \tilde{\phi} \rangle, \quad (18)$$

i.e. the condensate does not violate parity, we get the FESR constraints:

$$\sum_{i=e, e^*} \lambda_i^2 - \frac{N}{32\pi^2} t_c \simeq \langle \phi^+ \phi \rangle, \quad (19a)$$

$$\sum_{i=e, e^*} \lambda_i^2 m_i^2 - \frac{N}{64\pi^2} t_c^2 \simeq \frac{\alpha_H}{64\pi} \langle F^2 \rangle + \frac{3}{2} \left(\frac{N^2-1}{N^2} \right) \pi \alpha_H \left\{ 3 \langle \phi^+ \phi \rangle^2 - \frac{1}{3} \langle \phi \tilde{\phi} \rangle^2 \right\}. \quad (19b)$$

One can still simplify the result in Eq. (19b) by using the relation:

$$\langle \phi^+ \phi \rangle \simeq \frac{1}{\sqrt{2}} \langle \phi \tilde{\phi} \rangle \quad (20)$$

obtained from the analysis of the Goldstone sector in a SUSY-QCD theory^{19c)}. One can also neglect the $\alpha_H \langle F^2 \rangle$ condensate in a SUSY-QCD theory²³⁾. Then, we deduce, using the positivity of λ_e^2 and for $SU(2)_H$:

$$m_{e^*}^2 \gtrsim \frac{t_c}{2} \left\{ \frac{1 + 84\pi^3 \alpha_H \frac{\langle \phi^+ \phi \rangle^2}{t_c^2}}{1 + 16\pi^2 \frac{\langle \phi^+ \phi \rangle}{t_c}} \right\}. \quad (21)$$

One can see that solving Eq. (21) for $m_{e^*} \simeq 70$ GeV, $t_c \simeq 1$ TeV² and $\alpha_H \simeq 1$ does not give a real solution for $\langle \phi^+ \phi \rangle$, i.e. the assumption of having a e^* -mass at 70 GeV is unrealistic. If the W-boson is also composite, one can see in the SUSY case that the analogue of Eq. (11) (modulo obvious change in the coefficient) is not altered by the SUSY vacuum condensate. Then, the conclusion for the non-supersymmetric case continues to hold here, i.e., the value of the

vacuum condensate $\langle \phi^+ \phi \rangle$ should be less than M_W^2 . Equation (21) shows that the mass squared of the e^* , μ^* ... should be at least larger than half of the compositeness scale, and a SUSY composite model with elementary or composite W-bosons cannot favour a e^* mass setting at the W-scale.

4. CONCLUSION

We can summarize the analysis of this paper by the following:

- a) In a non-supersymmetric and a supersymmetric model where the fermions (quarks, lepton) and the W-bosons are both composite of scalar and spinor fields, the hope to have excited fermions at the W-scale cannot be realized if the fermion spectrum respects the duality picture. We found that the mass of the excited fermion should be at least 0.5 TeV.
- b) The only model which can admit a 70 GeV mass for the excited fermion is the non-supersymmetric model where the W-bosons are elementary and where the vacuum condensate of scalar should be larger than $(1 \text{ TeV})^2$.

We believe that our results apply to most classes of composite models based on hyperstrong QCD-like theories at the TeV scale.

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Fig. 1 Leading QHD contribution to the two-point function propagator:
a) perturbative; b) $\langle \phi^+ \phi \rangle$; c) $\alpha_H \langle F^2 \rangle$

Fig. 2 Analogous to Fig. 1 but for the two-point vector function

Fig. 3 Sample of contributions due to the Yukawa term $\lambda \phi \phi^+$ in the SUSY Lagrangian:
a) $\langle \phi^+ \phi \rangle^2$; b) $\langle \phi_i^+ \phi_i \rangle \langle \phi_j^+ \phi_j \rangle$; c, d) $g \langle \lambda \phi \phi^+ \rangle$
in addition to the contributions in Fig. 1.

Table caption

Different contributions to the two-point function of fermions (Eq. 13) and of vector (Eq. 14) currents in a supersymmetric QHD theory where the matter superfield mass is zero. The results for the non-supersymmetric theory can be deduced from the former by eliminating the contributions of all diagrams containing the $\lambda \phi \phi$ coupling and the four-point scalar interactions. One should also substitute $\langle \phi^+ \phi + \tilde{\phi}^+ \tilde{\phi} \rangle$ by $2\langle \phi^+ \phi \rangle$ and take $\langle \phi \tilde{\phi} + \tilde{\phi}^+ \phi^+ \rangle = 0$. In the case of the vector current, one should divide by a factor two all diagrams which do not contain a fermion loop.

TABLE

LEPTON TWO-POINT FUNCTION: SUSY CASE

$$S(q) \equiv \frac{1}{2} (1 - \gamma_5) \not{q} S(-q^2 \equiv \Phi^2)$$

$$\frac{1}{\epsilon} \equiv \frac{1}{\epsilon} + \log \frac{\Phi^2}{4\pi\gamma^2} + \gamma_E \quad \text{for } d = 4 + 2\epsilon$$

$$S(\Phi^2) = - \frac{N}{32\pi^2} \left(\frac{1}{\epsilon} - 2 \right)$$

$$S(\Phi^2) = \frac{1}{2} \langle \phi^\dagger \phi + \tilde{\phi}^\dagger \tilde{\phi} \rangle \frac{1}{\Phi^2}$$

$$S(\Phi^2) = - \frac{1}{2\sqrt{2}} g_H \langle \phi^\dagger \lambda \psi + \tilde{\phi}^\dagger \bar{\lambda} \bar{\psi} \rangle \frac{1}{\Phi^4}$$

$$S(\Phi^2) = - \frac{1}{4\sqrt{2}} g_H \langle \phi^\dagger \lambda \psi + \tilde{\phi}^\dagger \bar{\lambda} \bar{\psi} \rangle \frac{1}{\Phi^4}$$

$$S(\Phi^2) = - \frac{1}{64\pi} \alpha_H \langle F^2 \rangle \frac{1}{\Phi^4}$$

$$S(\Phi^2) = \left(\frac{N^2 - 1}{N^2} \right) \frac{1}{4} \pi \alpha_H \left\{ \langle \phi^\dagger \phi + \tilde{\phi}^\dagger \tilde{\phi} \rangle^2 - \langle \phi \tilde{\phi} + \phi^\dagger \tilde{\phi}^\dagger \rangle^2 \right\} \frac{1}{\Phi^4}$$

$$S(\Phi^2) = \left(\frac{N^2 - 1}{N^2} \right) (-1) \pi \alpha_H \left\{ \langle \phi^\dagger \phi + \tilde{\phi}^\dagger \tilde{\phi} \rangle^2 \right\} \frac{1}{\Phi^4}$$

W TWO-POINT FUNCTION: SUSY CASE

$$\tilde{\pi}^{\mu\nu}(q) = \frac{N}{12\pi^2} \left(\frac{1}{\epsilon} - \frac{5}{3} \right) (q^2 g^{\mu\nu} - q^\mu q^\nu)$$

$$\pi^{\mu\nu}(q) = \frac{N}{24\pi^2} \left(\frac{1}{\epsilon} - \frac{8}{3} \right) (q^2 g^{\mu\nu} - q^\mu q^\nu)$$

$$\pi^{\mu\nu}(q) = -2 \langle \phi^\dagger \phi + \tilde{\phi}^\dagger \tilde{\phi} \rangle \frac{1}{q^2} (q^2 g^{\mu\nu} - q^\mu q^\nu)$$

$$\tilde{\pi}^{\mu\nu}(q) = -2\sqrt{2}g \langle \phi^\dagger \lambda \psi + \tilde{\phi}^\dagger \bar{\lambda} \bar{\psi} \rangle \frac{1}{q^4} q^\mu q^\nu$$

$$\tilde{\pi}^{\mu\nu}(q) = +2\sqrt{2}g \langle \phi^\dagger \lambda \psi + \tilde{\phi}^\dagger \bar{\lambda} \bar{\psi} \rangle \frac{1}{q^4} (q^2 g^{\mu\nu} - 2q^\mu q^\nu)$$

$$\tilde{\pi}^{\mu\nu}(q) = 4\sqrt{2}g \langle \phi^\dagger \lambda \psi + \tilde{\phi}^\dagger \bar{\lambda} \bar{\psi} \rangle \frac{1}{q^4} q^\mu q^\nu$$

$$\pi^{\mu\nu}(q) = -\frac{1}{16\pi} \alpha_H \langle F^2 \rangle \left(\frac{1}{\epsilon} + \frac{1}{2} \right) \frac{1}{q^4} (q^2 g^{\mu\nu} - q^\mu q^\nu)$$

$$\pi^{\mu\nu}(q) = -\frac{1}{24\pi} \alpha_H \langle F^2 \rangle \frac{1}{q^4} (q^2 g^{\mu\nu} + 2q^\mu q^\nu)$$

$$2b^2 \frac{\phi}{j} \left[2 \langle +\phi + \phi + \phi \phi \rangle + \right.$$

$$2 \langle \phi_+ \phi_- + \phi_- \phi_+ \rangle \left(\frac{2N}{1-2N} \right) \delta = (b)$$

$$- \frac{1}{2} \langle \phi + \phi + \phi + \phi \rangle \} \mu_B \left(\frac{2N}{2N} \right) \sigma_- = (b)$$

$$+ 2 \langle \phi_+ \phi_- + \phi_- \phi_+ \rangle \} n_{\mathcal{U}} \left(\frac{2N}{1-2N} \right) \sigma = \quad (6)$$

$$c_{b_n} b^{\frac{1}{2} \Phi} \{ z \langle \phi + \phi + \phi \phi \rangle + z \langle \phi + \phi + \phi + \phi \rangle \}^{\frac{1}{2} N} \mathcal{L} \left(\frac{zN}{1-zN} \right) \mathcal{E} = (b)_{c^{\frac{1}{2} N} \mathcal{L}}$$

$$+ \frac{1}{2} \langle \phi_+ \phi_+ + \phi_- \phi_- \rangle \left(\frac{2N}{1 - \frac{1}{2}N} \right)^{1/2} = (b)$$

$$(a b_m b_n - a_m b_n b) \frac{1}{f} \{ \phi + \phi + \phi + \phi \} -$$

$$2 \{ \phi + \phi + \phi + \phi \} \frac{1}{f} \left(\frac{2N}{f} \right) \phi - = (b)_{cm} \frac{1}{f}$$

$$(c b_a b - c a b_2 b) \frac{1}{4} \phi_{21} \frac{2\pi r}{1} = (b)_{ca} \frac{11}{1}$$

$$\left\{ a^6 b^6 \left(\frac{x}{2} - \frac{2}{x} \right) - a^4 b^2 b^4 \left(\frac{x}{2} - \frac{2}{x} \right) \right\}^{\frac{1}{4}} \left(\frac{1}{2} \right)^{\frac{1}{4}} = (6)^{\frac{1}{4}}$$

Fig. 1

Fig. 2

Fig. 3