The Anomalous Magnetic Moment of Quarks

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In the case of massless current quarks we find that the breaking of chiral symmetry usually triggers the generation of an anomalous magnetic moment for the quarks. We show that the kernel of the Ward identity for the vector vertex yields an important contribution. We compute the anomalous magnetic moment in several quark models. The results show that it is hard to escape a measurable anomalous magnetic moment for the quarks in the case of spontaneous chiral symmetry breaking.

Theoretically, the various hadronic electromagnetic form factors are usually described in terms of pole dominance together with contributions arising from virtual mesonic exchanges \[ \text{I} \]. A third contribution to the electromagnetic form factors should come from the quark microscopic interaction itself, in close analogy with QED. It is clear that these three scenarios should not be independent but just three different aspects of the same model. This desideratum can be achieved, at least qualitatively, in terms of a quark field theory displaying spontaneous breaking of chiral symmetry, \( S\chi SB \). In such a description any hadron, when seen from the trivial vacuum Fock space, appears as a collection of an infinite number of quark antiquark pairs together with the appropriate valence quarks. It happens that the contributions of this quark sea can be summarized in terms of a new set of valence quasiquarks which now carry the information on the details of the physical vacuum truncated a modified propagator \[ \mathcal{E} \]. In this fashion we recover the simplicity of the constituent quark picture. It is the role of the Ward identities to ensure charge conservation throughout this process. And this they do at the expenses of the quark identities to ensure charge conservation throughout this process. And this they do at the expenses of the quark identities to ensure charge conservation throughout this process. And this they do at the expenses of the quark identities to ensure charge conservation throughout this process.

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The Ward Identity, \( i\sigma_{\mu\nu}S(p+q/2)\Gamma^\mu S(p-q/2) = S(p+q/2) - S(p-q/2) \)

\[ \Leftrightarrow q_\mu \Gamma^\mu = IS^{-1}(p+q/2) - iS^{-1}(p-q/2) \]

is obeyed both by the bare vertex \( \Gamma_0^\mu \) and by the Bethe-Salpeter vertex \( \Gamma^\mu \). We will show that in the limit of...
small momentum $q$, this identity has the following solution for the vertex,

$$\Gamma^\mu(p, q) = i \frac{\partial}{\partial p_\mu} S^{-1}(p) + q_\nu T^\nu\mu(p) + o(q^2)$$

(5)

where $q_\nu T^\nu\mu(p)$ is defined as the kernel which is not determined by the Ward identity,

$$q_\nu [q_\rho T^\rho\nu(p)] = 0 .$$

(6)

The Ward identity ensures that charge conservation survives renormalization. However it does not constrain the kernel, which is a signature of the renormalization. In particular the kernel contributes to the anomalous magnetic moment are respectively, the self energy and the kernel to the anomalous magnetic moment are both infrared divergent but their sum is finite:

$$S(\pi)$$

(7)

The vertex is given by,

$$\Gamma^\mu = \Gamma^\mu_0 - i \partial^\mu \Sigma + q_\nu T^\nu\mu$$

(8)

and, up to first order in $\alpha$, the contributions from the self energy and the kernel to the anomalous magnetic moment are respectively,

$$\left(-1 + \frac{2}{\pi} \ln \frac{\lambda}{\Lambda}\right) \frac{\alpha}{2\pi}, \quad \left(2 + 2 \ln \frac{\lambda}{\Lambda}\right) \frac{\alpha}{2\pi}$$

(9)

In the case of actual QED, where $\lambda \to 0, \Lambda \to \infty$, they are both infrared divergent but their sum is finite: $\alpha/2\pi$.

As for QCD, there has been a considerable effort on how to derive quark models by integrating out, under various approximations, the gluonic degrees of freedom. An interesting and promising approach is provided by the cumulant expansion of the interaction term of the QCD Lagrangian [3]. A non-local Nambu Jona-Lasinio type Lagrangian (NJL) is obtained when we retain only bilocal correlators. Therefore we hold the view that such quark models are appropriate to study electromagnetic properties of hadrons, even for light quarks, provided we have small enough photon momenta and the physics of chiral symmetry breaking is treated correctly. Therefore, at this stage, rather than focusing on a specific example of NJL we will study the static electromagnetic properties of a wide class of quark effective quartic interactions.

In quark models with dynamical $S\chi SB$, the vertex $\Gamma^\mu$ is a solution of the Bethe-Salpeter equation,

$$\Gamma^\mu(p, q) = \Gamma^\mu_0 - i \int \frac{d^4 p'}{(2\pi)^4} V(p' - p, p' + p, q) \Omega_a S(p'_1)$$

$$\Gamma^\mu(p', \frac{q}{2}) S(p'_2) \Omega_a - V(q, p' + p, -p' + p) \Omega_a$$

$$\text{tr} \{S(p'_1) \Gamma^\nu(p', q) S(p'_2) \Omega_a\}$$

(11)

where the $-1$ factor from the fermion loop was included in the tadpole term. The momentum dependence of the potential is only assumed to conserve the total momentum, and in this case it depends on 3 momenta. The Dirac, flavor and color structure of the interaction is determined by the $\Omega_a$ matrices. In order to have dynamical $S\chi SB$, we require this structure to be chiral invariant. Substituting the Ward Identity in the ladder Bethe Salpeter equation for the vertex we get,

$$\int d^4 p' \frac{d^4}{(2\pi)^4} V(p' - p, p' + p, q) \Omega_a [S(p'_1) - S(p'_2)] \Omega_a$$

$$- V(q, p' + p, -p' + p) \Omega_a \text{tr} \{S(p'_1) - S(p'_2) \Omega_a\}$$

(12)

For particular cases of the potential $V(p'_1 - p_1, p'_1 + p_2, p'_1 - p'_2)$ we recover the BCS mass gap equation,

$$S^{-1}(p) = S_0^{-1}(p) - \frac{p - p'}{p'}$$

(13)

provided that either the rainbow diagram vanishes or,

$$V(p' - p, p' + p, q) = V(p' - p, p' + p, 0) ,$$

(14)

and that either the tadpole diagram vanishes or,

$$V(q, p' + p, -p' + p) = V(0, p' + p, -p' + p) .$$

(15)

Equation [13] can be written,

$$S^{-1}(p) = S_0^{-1}(p) - \int \frac{d^4 p'}{(2\pi)^4} V(p' - p, p' + p, 0) \Omega_a S(p'_1) \Omega_a$$

$$- V(0, p' + p, -p' + p) \Omega_a \text{tr} \{S(p'_1) \Omega_a\}$$

(16)

Now we insert the expression [1] for $\Gamma^\mu$ in the Bethe Salpeter equation [1], in order to find the kernel $qT$ and expand it up to first order in $q$. The equation for the tensor $T$, which is antisymmetric, is then,

$$\Gamma^{\nu\mu} = T_0^{\nu\mu} - i \int \frac{d^4 p}{(2\pi)^4} V \Omega_a (1 - tr) \{ST^{\nu\mu} S \Omega_a\}$$

$$T_0^{\nu\mu} - \frac{1}{2} \int \frac{d^4 p}{(2\pi)^4} V \Omega_a (1 - tr) \{T^{\nu\mu} S \Omega_a\}$$

$$J^{\nu\mu} = \partial^\nu (S^{-1} \partial^\mu (S - \partial^\mu (S) S^{-1} \partial^\nu (S))$$

(17)
This is a self consistent forced linear integral equation. Let us consider a general quark propagator, solution of the mass gap equation, of the form,

$$ S(p^\mu) = \frac{iF(p)}{\not p - M(p)} $$

(18)

where $p = \sqrt{p^\mu p_\mu}$. The integrand $J^{\mu\nu}$ is then,

$$ J^{\mu\nu} = -i \frac{F}{(p^2 - M^2)^2} \left[ \frac{1}{2} \{ \not p, \{ \gamma^\nu, \gamma^\mu \} \} + M[\gamma^\nu, \gamma^\mu] \\
- \frac{M}{p} (p^\nu [\not p, \gamma^\mu] - p^\mu [\not p, \gamma^\nu]) \right] $$

(19)

where the dot superscript denotes $d/dp$. In general we find,

$$ T^{\nu\mu} = t_1(p) \{ \not p, \{ \gamma^\nu, \gamma^\mu \} \} + t_2(p) M[\gamma^\nu, \gamma^\mu] \\
+ t_3(p)(p^\nu [\not p, \gamma^\mu] - p^\mu [\not p, \gamma^\nu]) $$

(20)

Up to $o(q^2)$ the electromagnetic current of the quark is then,

$$ j^\mu = \frac{e}{F} \bar u \gamma^\mu p^\mu M - (\not p - M) \frac{p^\mu F}{p F} + Fq^\nu T^{\nu\mu} u \\
= \frac{e(1 - \lambda F)}{F} \bar u \left[ \gamma^\mu + a \left( \frac{i \sigma^\mu_\nu}{2M} q^\nu \right) \right] u, \\
a = \frac{M + 4M^2 F(2t_1 + t_2)}{1 - M} $$

(21)

where the mass shell condition $p = M$ was used together with the Gordon identities. The anomalous magnetic moment $a$ turns out to be independent of $t_3$ and $F$. However the dependence on $M$ is crucial in models where $t_1$ and $t_2$ are finite. In those models $a$ can be thought as a measure of $S\chi SB$. The quark condensate $\langle \bar qq \rangle$ is also a functional of the dynamically generated mass,

$$ \langle \bar qq \rangle = - \quad \quad = n_c \operatorname{tr} \int \frac{dp}{(2\pi)^4} S(p) $$

(22)

where the trace sums colors with $n_c = 3$, but the flavor is kept fixed. Thus, at the onset of the spontaneous $\chi$SB, we will obtain an implicit relation between $a$, $\langle \bar qq \rangle$ and the constituent quark mass, which were simultaneously vanishing before the occurrence of this phase transition and now become non-zero.

We will now compute $F$, $M$, $a$ and $\langle \bar qq \rangle$ in particular models which are paradigmatic cases of chiral symmetry breaking and comply with the constraints of the Ward identity.

Model I is the first original NJL model [8]. The Lagrangian of model I is,

$$ \mathcal{L}_I = \bar q i \not \! D q + G \left( (\bar qq)^2 - (\bar q \gamma_5 q)^2 \right) $$

(23)

where $\mathcal{L}_I$ is specific to the case of 1 flavor, but its results are similar to the ones of flavor symmetric $U_A(n_f)$ extended NJL models. The equations will be solved in the momentum representation. As usual the integrals are done in Euclidean space. A momentum cutoff $\Lambda$ is included in order that the integral in the loop momentum is finite. Since the cutoff cannot be ascribed to the potential which has to be constant in momentum space, must be included in the propagator,

$$ S(p) = \frac{i F(p)}{\not p - M + i\epsilon}, \quad F(p) \to \Theta_{\text{Euclidean}}(\Lambda - p). $$

(24)

With a constant potential and this momentum cutoff, the loops turn out to be constant, independent of the external momentum $p$. It is convenient to evaluate the integrals,

$$ I_1 = - \int \frac{d^4p}{(2\pi)^4} \frac{i F}{(p^2 - M^2)^2} $$

$$ = \frac{1}{16\pi^2} \left[ \frac{A^2 - M^2}{\Lambda^2 - M^2} \ln \left(1 + \frac{\Lambda^2}{M^2}\right) \right], $$

$$ I_2 = i \int \frac{d^4p}{(2\pi)^4} \frac{F}{(p^2 - M^2)^2} $$

$$ = \frac{1}{16\pi^2} \left[ - \frac{A^2}{\Lambda^2 + M^2} + \ln \left(1 + \frac{\Lambda^2}{M^2}\right) \right], $$

(25)

where the solid angle $2\pi^2$ is included. The mass gap equation is,

$$ \hat p - M = \hat p - 2G \int \frac{d^4p}{(2\pi)^4} \frac{i F}{(p^2 - M^2)^2} \left[ (\hat p + M) \\
- \gamma_5 (\hat p + M) \gamma_5 - tr \{ \hat p + M \} \right] $$

(26)

With the solutions $M = 0$ or $1 = 8 n_c G I_1(M, \Lambda)$. The parameters $\Lambda$ and $G$ are determined once the quark dynamical mass and the quark condensate are fixed. We now study the kernel in model I. Because the integrals are constant, the antisymmetric tensor $T$ is independent of $p$. Thus $T^{\nu\mu}$ has to be of the $t_2$ type, proportional to $[\gamma^\nu, \gamma^\mu]$. Including the structure factors $\Omega_q$ we find that the tadpole-like term vanishes since $\sigma^{\mu\nu}$ and $\sigma^{\mu\nu} \gamma_5$ have a null trace. In this case of model I the rainbow diagram also cancels since the structure factor $1 \otimes 1 - \gamma_5 \otimes \gamma_5$ projects on the terms with an odd number of Dirac $\gamma$ matrices, of type $t_1$ but $[\gamma^\nu, \gamma^\mu]$ is even. Thus model I produces no kernel for the vector vertex and no anomalous magnetic moment for the quark $\bar q$.

Model II is the second original NJL model [8]. The Lagrangian is,

$$ \mathcal{L}_{II} = \bar q i \not \! D q + G \left( (\bar qq)^2 - (\bar q \gamma_5 \tau q)^2 \right) $$

(27)

where $\mathcal{L}_{II}$ is used for 2 flavors $v$ and $d$. It only has an $SU(2)_A$ symmetry and breaks $U(1)_A$ from the onset. Its
results are similar to those of flavor symmetric $SU_A(n_f)$ extended NJL models. The anzats for the propagator is that of Eq. [3], and model II only differs from model I in the algebra. The mass gap equation is changed since $\vec{\tau} \cdot \vec{\tau} = 3$ in the fermion line. We get,

$$\hat{p} - M = \hat{p} - 2G \int \frac{d^4p}{(2\pi)^4} \frac{iF}{p^2 - M^2} \left[ (\hat{p} + M) \right]$$

$$-2n_f^2 - \frac{1}{n_f} \gamma_5 (\hat{p} + M) \gamma_5 - t \{ \hat{p} + M \}$$

$$\Rightarrow M = 0 \text{ or } 1 = 2 \left( 2n_f^2 - \frac{1}{n_f} - 1 + 4n_f n_c \right) G_1(M, \Lambda).$$

were $n_f$ and $n_c$ stand respectively for the number of flavors and colors. The rainbow diagram contributes in this case. The tadpole diagram contribution also changes. The preferred values for the parameters $\Lambda$ and $G$ are $\Lambda = 1.65$GeV and $G = 1.23$GeV$^{-2}$ which yield $M = 333$GeV, $\langle \bar{q} \gamma u \rangle = -0.25$GeV$^3$ and $f_\pi = .09$GeV. As in model I, the tadpole will not contribute to the antisymmetric tensor which will be again of the $t_2$ type, $T^{\nu\mu} = t_2 [\gamma^\nu \gamma^\mu]$. The first order term is a function of,

$$\int \frac{d^4p}{(2\pi)^4} J^{\nu\mu} = -I_2 \ M [\gamma^\nu, \gamma^\mu].$$

In order to evaluate the higher order terms, we calculate,

$$\int \frac{d^4p}{(2\pi)^4} S \ M [\gamma^\nu, \gamma^\mu] S = -iM^2 I_2 \ q_\nu \ M [\gamma^\nu, \gamma^\mu]$$

In this case we have two flavors with two different charges $e_u = \frac{2}{3}$, $e_d = -\frac{1}{3}$, and two anomalous magnetic moments $a_f$,

$$e_u I_2 = G I_2 \left[ 0 \left( \frac{e_u}{2} - M^2 e_u \right) + (-2) \left( \frac{e_d}{2} - M^2 e_d t_d \right) \right]$$

$$\text{ (u} \leftrightarrow \text{ d) }$$

The natural parameter is $2G M^2 I_2 (\Lambda, M) = 0.004$. Inverting this equation we find the solution,

$$a_u \simeq -2(2GM^2 I_2) \frac{e_d}{e_u} = 0.004 \Rightarrow M_u = 339 \text{ MeV}$$

$$a_d \simeq -2(2GM^2 I_2) \frac{e_u}{e_d} = 0.016 \Rightarrow M_d = 327 \text{ MeV}$$

Although this effect is small, it has the right sign to correct the $M_u$ and $M_d$ inversion. If the tadpole term was removed from the mass gap equation then the $a_d - a_u$ would be bigger. This is possible for instance when the potential has a $\vec{\lambda} \cdot \vec{\lambda}$ dependence, being $\lambda$ the Gell-Mann matrices.

Model III is the simplest QCD inspired model. The Lagrangian is,

$$L_{III} = \bar{q} \gamma \hat{p} + \frac{1}{2} \bar{q}(x) \gamma_5 \vec{\lambda} \cdot \vec{q}(x) \int d^4y V(x-y) \bar{q}(y) \gamma_5 \vec{\lambda} \cdot \vec{q}(y)$$

In the case of model III, the Dirac structure $\gamma^\mu \otimes \gamma_\mu$ is $U_A(n_f)$ chiral invariant. For $V(p)$ we will choose a color confining square well potential because of its calculational simplicity.

$$V(p' - p) = -G \Theta_{\text{Euclidian}} (\Lambda - |p' - p|),$$

The mass gap equation is,

$$\hat{p} - M = \frac{4i}{3} \int \frac{d^4p'}{(2\pi)^4} V(p' - p) \frac{(\hat{p}' + 4M)}{p'^2 - M^2}$$

which includes the color factor of $\frac{4}{3}$. We now calculate the kernel. The first order term for the kernel is a functional of,

$$\gamma^\alpha J^{\nu\mu} \gamma_\alpha = iF \{ \hat{p}, [\gamma^\nu, \gamma^\mu] \}$$

The terms with two gamma matrices, of the form $\sigma^{\mu\nu}$ are now cancelled by the $\gamma^\alpha \otimes \gamma_\alpha$ of the interaction, and only the $t_1$ type term remains. Therefore this model differs from the previous ones insofar it covers the form factor $t_1$. The self consistent equation for the antisymmetric tensor $T$ will also close,

$$\gamma^\alpha S \{ \hat{p}, [\gamma^\nu, \gamma^\mu] \} S \gamma_\alpha = 2F^2 \frac{p^2 + M^2}{(p^2 - M^2)^2} \{ \hat{p}, [\gamma^\nu, \gamma^\mu] \}$$

We get,

$$T^{\nu\mu} = t \{ \hat{p}, [\gamma^\nu, \gamma^\mu] \}$$

$$t(p) = t_0 - \frac{8}{3} \int \frac{d^4p'}{(2\pi)^4} V(p' - p) \frac{p'}{p^2} F^2(p')$$

$$\frac{p^2 + M^2}{(p^2 - M^2)^2} t(p'),$$

$$t_0(p) = -\frac{2}{3} \int \frac{d^4p'}{(2\pi)^4} V(p' - p) \frac{p'}{p^2} F(p')$$

For the euclidean integration it is convenient to evaluate the angular integrals,

$$I_3(\ p', p) = G \int_{-1}^{+1} d\theta \left( \Lambda - \sqrt{p'^2 + p^2 - 2wp'p} \right)$$

$$= G(1 + I_6) \theta(1 - I_6) (1 + I_6) + 2G \theta(I_6 - 1) > 0$$

$$I_4(\ p', p) = G \int_{-1}^{+1} d\theta \left( \Lambda - \sqrt{p'^2 + p^2 - 2wp'p} \right)$$

$$= G \frac{I_3^2}{2} \theta(1 - I_6) (1 + I_6) > 0$$

$$I_5 = \frac{\Lambda^2 - p'^2 - p^2}{2wp'p}$$

and the nonlinear integral mass gap equation for $F$ and $M$, the integral for $t_0$, the linear integral equation for $T$, and the integral for $\langle \bar{q}q \rangle$ can be solved simultaneously,
\[ F(p) = \left[ 1 + \int_0^\infty dp \frac{I(p',p)}{6\pi^2p} \frac{p'^4}{p'^2 + M^2(p')} F(p') \right]^{-1} \]

\[ M(p) = F(p) \int_0^\infty dp \frac{I(p',p)}{3\pi^2} \frac{p'^3 F(p')}{p'^2 + M^2(p')} M(p') \]

\[ t_o(p) = \int_0^\infty dp \frac{I(p',p)}{24\pi^2p} \frac{p'_4 F(p')}{[p'^2 + M^2(p')]} \]

\[ t(p) = t_o - \int_0^\infty dp \frac{I(p',p)}{6\pi^2p} \frac{p'^4 F(p')}{[p'^2 + M^2(p')]} t(p') \]

\[ \langle \bar{q} q \rangle = -\int_0^\infty dp \frac{3}{2\pi^2} \frac{p'^4 F(p) M(p)}{p'^2 + M^2(p)} \]

(37)

The mass term has a trivial solution \( M = 0 \) and another solution which breaks spontaneously chiral symmetry. A dimensional simplification occurs if we work in units of \( \Lambda = 1 \). In this case the only parameter is \( G \) which is now adimensional. We find a critical value \( G_c = 132 \) above which chiral symmetry occurs. In Fig. 1 we depict the values of \( M, \langle \bar{q} q \rangle \) and \( a \). We solve the integral equations numerically for \( F, M \) and \( t \) with the Gaussian iterative method and using the Gaussian integration [10]. We find that at \( p^2 = -1 \) these functions decrease by a factor of just \( .9 \to .7 \). Since we cannot continue analytically the numerical solution we use the approximation of nearly constant \( F, M \) and \( t \) and compute the mass and the anomalous magnetic moment for \( p = 0 \). The literature prefers \( a < \langle \bar{q} q \rangle = -(0.25 \text{ GeV})^3 \). A dynamical quark mass \( M = 0.33 \text{ GeV} \) would correspond to \( G = 245 \Lambda^{-2}, \Lambda = 0.74 \text{ GeV} \) and \( a = 0.15 \). If we now consider a \( M = (1 + a)0.33 \text{ GeV} \) then the lowest possible condensate is \( \langle \bar{q} q \rangle = -(0.28 \text{ GeV})^3 \) which corresponds to \( G = 300 \Lambda^{-2}, \Lambda = 0.69 \text{ GeV} \), \( M = 0.42 \text{ GeV} \) and \( a = 0.28 \), see Fig. 1.

NJL models I and II are the simplest models with chiral symmetry breaking. In the NJL model I the anomalous magnetic moment \( a \) vanishes. In the model II the \( U(1) \) breaking interaction yields a too small \( a \), which nevertheless provides an example of an isospin dependence for \( a \) and, therefore, contributes to the \( u - d \) mass inversion. The reason for the smallness of the anomalous magnetic moment stems from the presence of tadpole contributions and were not for this contribution and we would have obtained a much larger \( a \). This is precisely the case of model III where a larger \( a \) is derived, compatible with the nonrelativistic constituent quark models. We also find that \( M, a, \langle \bar{q} q \rangle \), are functions of \( G - G_c \) with critical exponents which are respectively 1, 2 and 1. The present work constitutes a first step on a more elaborate model unifying hadronic spectroscopy (including decay widths) and the electromagnetic form factors which have been shown to be consistent with the simple quark constituent picture precisely because of \( S \chi SB \).

REFERENCES

[1] H. Ito, Phys. Rev C 52, R1750, (1995); H. Forkel, M. Nielsen, X Jin, T. Cohen, Phys Rev C 50, 3108, (1994).
[2] A. Le Yaouanc, L.Oliver, O. Pene and J-C. Raynal, Phys Rev D 29, 1233 (1984); A. Le Yaouanc, L.Oliver, S. Ono, O. Pene and J-C. Raynal, ibid 31, 137 (1985); P.Bicudo and J. Ribeiro, Phys Rev D 42, 1611 (1990).
[3] P.Bicudo and J.Ribeiro Phys Rev D 42, 1635 (1990); P.Bicudo and J.Ribeiro Phys Rev C 55, 834 (1997).
[4] N. Isgur, G. Karl, Phys. Rev. D 18, 4178.
[5] S. Adler and A. C. Davis, Nucl. Phys. B 224, 469 (1984).
[6] H. G. Dosh, Phys Lett. B190 177 (1987); H. G. Dosh and U. Marquard, Nuc Phys. A560 333 (1993); N. Brambilla and A. Vairo Phys. Lett B407 167 (1997); Yu. A. Simonov [hep-ph/9712248].
[7] Y. Nambu, G. Jona-Lasinio, Phys. Rev. 122, 345 (1961).
[8] U. Vogl, M. Lutz, S. Klimt and W. Weise, Nucl. Phys. A 516, 469 (1990); J. Singh, Phys. Rev. D 31, 1097 (1985).
[9] Y. Nambo, G. Jona-Lasinio, Phys. Rev. 124, 246 (1961); V. Bernard, Phys. Rev. D 34, 1601 (1986).
[10] Y. Dai, Z. Huang and D. Liu, Phys. Rev. D 43, 1717 (1991).

FIG. 1. Functions of the quark dynamical mass \( M \), anomalous magnetic moment \( a \) and quark condensate as functions of the adimensional coupling \( G \) for model III.