The $(0^+, 1^+)$ heavy meson multiplet in an extended NJL model

D. Ebert$^1$, T. Feldmann$^1$,  
*Institut für Physik, Humboldt–Universität,  
Invalidenstrasse 110, D–10115 Berlin, Germany

R. Friedrich$^{2,3,4}$, H. Reinhardt$^2$,  
*Institut für Theoretische Physik, Universität Tübingen,  
Auf der Morgenstelle 14, D–72076 Tübingen, Germany

Abstract

In this letter we reconsider the previously given description of heavy mesons within a bosonized extended NJL model that combines heavy quark and chiral symmetry. In that work the naive gradient expansion of the quark determinant was used, which satisfactorily works in the light sector but does not adequately describe the heavy $(0^+, 1^+)$ mesons. By investigating the exact momentum dependence of the quark loop we demonstrate that the naive gradient expansion in the heavy sector is not the right method to treat the unphysical $qar{q}$–thresholds which would be absent in confining theories. We propose a modified gradient expansion which adequately extrapolates from the low–momentum region beyond threshold. This expansion gives a satisfactory description even of the $(0^+, 1^+)$ heavy mesons whose masses are significantly above threshold.

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$^4$e–mail: friedric@ptdec5.tphys.physik.uni-tuebingen.de
1 Introduction

The observation of new symmetries for infinitely heavy quarks in QCD together with the formulation of the Heavy Quark Effective Theory (HQET) [1] has led to a number of theoretical and phenomenological investigations. One interesting application is to exploit these symmetries in the formulation of tractable quark models [2, 3, 4].

In ref. [4] we have combined chiral symmetry of light quarks with heavy quark spin and flavor symmetry in an extension of the Nambu–Jona–Lasinio (NJL) model. In order to be as close as possible to former investigations in the light sector [5, 6], we have fixed most of the parameters from light meson physics. For evaluating the quark determinant we have employed a linear gradient expansion.

While our numerical investigations, relating heavy meson properties and model parameters, have been successful for the \((0^-, 1^-)\) meson doublet of heavy quark spin symmetry, our approach seems to fail for the heavier \((0^+, 1^+)\) states of opposite parity. This seems somewhat surprising since such a problem does not arise in the analogous case of the \(\rho\) or \(a_1\) mesons in the light sector.

The aim of this work is to clarify this problem by a re-investigation of the gradient expansion in both the heavy and the light meson sector. The gradient expansion (or equivalently expansion of quark loops in external momenta) is usually performed, in order to avoid unphysical quark thresholds in NJL–type models without quark confinement. In fact, assuming such models valid in the low–momentum region far below the thresholds, a gradient expansion can afterwards be interpolated to external momenta of even a few 100 MeV above the threshold where it still predicts satisfactory results [3, 4].

Nevertheless, the situation for heavy and light mesons is qualitatively different due to the different analytical structure of quark loop expressions involving a heavy quark propagator \((v \cdot k + i\epsilon)^{-1}\). Therefore, one should look carefully how to define the interpolation procedure to larger external momenta which is necessary to describe the heavier states.

The organization of the letter is as follows: In section 2 we will give a short review of the definition of the extended NJL model for heavy quark flavors and of the result for the heavy meson self–energy which defines masses and renormalization factors of heavy mesons [4]. We show that the \((0^-, 1^-)\) and \((0^+, 1^+)\) heavy mesons cannot be described simultaneously by one set of parameters. The next section 3 is devoted to a detailed analysis of the gradient expansion by evaluating the exact dependence on external momenta of the pertinent loop diagrams in both the light and heavy sector. Our analysis shows that a rather different way of interpolation is required for the heavy meson self–energy as compared to the light sector. This suggests a modified gradient expansion which is valid in both the light and heavy sector. The modified gradient expansion is shown to lead to the desired improvement in numerical results. Finally, some concluding remarks are given in section 4.

2 Free effective meson lagrangian and lowest–order gradient expansion

In ref. [4] we have presented an extension of the NJL model which combines chiral symmetry for light quarks with heavy quark symmetries for heavy quark fields defined by

\[
Q_v(x) = \frac{1 + \frac{g}{2} \exp(imQv \cdot x)Q(x)}{2}.
\]

We do not want to give technical details of the employed bosonization procedure but simply quote the result for the effective meson lagrangian when all the quark fields have been integrated.
\[ \mathcal{L} = -iN_c \text{Tr} \ln iD - \frac{1}{4G_1} \text{tr}_F \left[ \Sigma^2 - \tilde{m}_0 (\xi \Sigma \xi + \xi^\dagger \Sigma^\dagger) \right] + \frac{1}{4G_2} \text{tr}_F \left[ (V_\mu - V_\pi^\mu)^2 + (A_\mu - A_\pi^\mu)^2 \right] + \frac{1}{2G_3} \text{Tr} \left[ (\overline{H} + \overline{K}) (H - K) \right], \]  \hspace{1cm} (2)

where \( N_c = 3 \) is the color factor, \( \tilde{m}_0 \) is the current mass matrix of light quark flavors and \( G_1, G_2, G_3 \) are coupling constants of four–quark interactions. Furthermore, 

\[ \frac{1}{iD} = \frac{1}{i(-\nabla + \Sigma + V + A + i\gamma_5 \nabla) - (\overline{H} + \overline{K}) (iv \cdot \partial)^{-1} (H + K)} \]  \hspace{1cm} (3)

is the Dirac operator for the light quarks, which contains the several light \((\Sigma, V_\mu, A_\mu)\) and heavy meson fields \((H, K)\). In the heavy quark limit, heavy meson fields are organized in spin symmetry doublets

\[ H = P_+ (i\Phi^5 \gamma_5 + \Phi^\mu \gamma_\mu) \hspace{0.5cm}, \hspace{0.5cm} v_\mu \Phi^\mu = 0 \hspace{0.5cm}, \]  \hspace{1cm} (4)

\[ K = P_+ (\Phi + i\Phi^5 \gamma_\mu \gamma_5) \hspace{0.5cm}, \hspace{0.5cm} v_\mu \Phi^5 = 0 \hspace{0.5cm}, \]  \hspace{1cm} (5)

with \( \Phi \) being the heavy scalar, \( \Phi^5 \) the heavy pseudoscalar, \( \Phi^\mu \) the heavy vector and \( \Phi^5 \) the heavy axial vector field, respectively. \( V, A \) denote the light vector and axial–vector fields. For the light octet of Goldstone bosons we use the common non–linear representation \( \xi = \exp(i\pi/F) \) where \( \pi = \pi^a \lambda^a / 2 \) and \( F \) is the bare decay constant. Vector and axial–vector expressions are built via \( V^\mu_\pi = i/2(\xi \partial_\mu \xi^\dagger + \xi^\dagger \partial_\mu \xi) \), \( A^\mu_\pi = i/2(\xi \partial_\mu \xi^\dagger - \xi^\dagger \partial_\mu \xi) \). Finally, the light scalar field \( \Sigma \) achieves a non–vanishing vacuum expectation value indicating the spontaneously breaking of chiral symmetry.

From (3) we have derived self–energy expressions for mesons as well as interaction terms between heavy and light mesons. This is achieved by evaluating \( \text{Tr} \ln iD \) in terms of quark loops that are regularized by Schwinger’s gauge invariant proper–time method. Note that due to chiral and heavy quark symmetry, the heavy sector is controlled by only one coupling constant \( G_3 \) for both \( H \) and \( K \) fields, while all other parameters are fixed from the light sector. The self–energy term for heavy mesons \( \Pi_{H,K} \) as a function of the external momentum \((v \cdot p)\) is represented in figure [3] and calculated as

\[ -\text{tr}_D \left[ \overline{H} \Pi_{H}(v \cdot p) H \right] + \text{tr}_D \left[ \overline{K} \Pi_{K}(v \cdot p) K \right] = iN_c \int_{\text{reg}} d^4k \text{tr}_D \left[ \frac{\hat{k} - \hat{p} + m^i}{((k - p)^2 - (m^i)^2)(v \cdot k + i\epsilon)} (\overline{H}^i + \overline{K}^i) \right], \]  \hspace{1cm} (6)
where \( i = u, d, s \) is a light flavor index. Remember that the heavy \((Q_q \bar{u})\)–meson lagrangian in the heavy quark limit is given by

\[
\mathcal{L}_0^{\text{heavy}} = -\text{tr}_D \left[ \mathbf{T}^i (iv \cdot \partial - \Delta M_H) H^i \right] + \text{tr}_D \left[ \overline{K}^i (iv \cdot \partial - \Delta M_K) K^i \right],
\]

(7)

where \( \Delta M_{H,K}^i = M_{H,K}^i - m_Q^i \), with \( M_{H,K}^i \) being the heavy meson mass. Let us compare the inverse propagator term \((v \cdot p - \Delta M_{H,K}^i)\) resulting from (7) with the self–energy expression \( \Pi_{H,K}^i \) of (6) and the term \( \sim 1/G_3 \) in (2). First observe that \( \Pi_{H,K}^i \) is a function of \( v \cdot p \) only. Expanding \( \Pi_{H,K}^i \) around the on–shell value \( v \cdot p = \Delta M_{H,K}^i \), one obtains

\[
\Pi_{H,K}^i (v \cdot p) = \Pi_{H,K}^i (\Delta M_{H,K}^i) + \Pi_{H,K}^i (\Delta M_{H,K}^i) \left( v \cdot p - \Delta M_{H,K}^i \right) + O((v \cdot p - \Delta M_{H,K}^i)^2).
\]

(8)

With this expansion one reads off from (6) that the meson mass is determined by

\[
I_3^i (\Delta M_{H,K}^i) \left( \Delta M_{H,K}^i \pm m_i \right) + I_1^i - \frac{1}{2G_3} = 0
\]

(9)

and the \( Z \)–factors for the necessary field renormalization \( H^i, K^i \rightarrow (Z_{H,K}^i)^{-1/2} H^i, K^i \) are given by

\[
Z_{H,K}^i (\Delta M_{H,K}^i) = \left( I_3^i (\Delta M_{H,K}^i) + 2I_3^i (\Delta M_{H,K}^i) (\Delta M_{H,K}^i \pm m_i) \right)^{-1},
\]

(10)

where \( I_1^i, I_3^i \) are loop integrals defined by

\[
I_1^i = \frac{iN_c}{16\pi^4} \int_{\text{reg}} \frac{d^4k}{k^2 - (m_i)^2} = \frac{N_c}{16\pi^2} (m_i)^2 \Gamma(-1, (m_i)^2/\Lambda^2),
\]

(11)

\[
I_3^i (v \cdot p) = -\frac{iN_c}{16\pi^4} \int_{\text{reg}} \frac{d^4k}{(k^2 - (m_i)^2)(v \cdot k + v \cdot p + i\epsilon)} \frac{1}{(v \cdot p)^2} \Gamma \left( 0, (m_i)^2 - x(v \cdot p)^2 / \Lambda^2 \right)
\]

\[
+ \sqrt{\pi} \sqrt{(m_i)^2 - (v \cdot p)^2} \Gamma \left( -\frac{1}{2}, (m_i)^2 - (v \cdot p)^2 / \Lambda^2 \right).
\]

(12)

\( I_3^i (v \cdot p) = 1/2 \, dI_3^i (v \cdot p) / d(v \cdot p) \) and \( \Gamma(\alpha, x) \) is the incomplete gamma function.

As a consequence of the heavy flavor symmetry the solutions of (6) do not depend on the heavy quark mass. But they do depend on the light quark constituent masses and so are influenced by the explicit and spontaneous chiral symmetry breaking. For that reason we have supplied the mass shifts \( \Delta M_{H,K}^i \) with the index \( i \) of the light quark flavors.

It is obvious that for \( v \cdot p > m_i \) the quark loop produces an imaginary part due to the unphysical quark–antiquark threshold. Clearly, one must neglect this imaginary part of the meson self energy, since it would yield a nonvanishing decay width \( \Gamma(H^i, K^i \rightarrow \bar{q}^i + Q) \) forbidden by confinement. Its formal appearance is obviously related to the fact that confinement is not taken into account in our model explicitly. The simplest way to avoid these unphysical two–quark thresholds is to expand \( \Pi_{H,K}^i \) to first order around \( v \cdot p = 0 \) in analogy to the gradient

\textsuperscript{5} Note that the heavy meson field has the mass dimension 3/2.
expansion of the light meson sector, where one expands to first order around $p^2 = 0$ \cite{5,6}. This lowest–order expansion in $(v \cdot p)$ has been performed in ref. \cite{4}, leading to

$$
\Delta M_{H,K} = \left( I_3^i(0) \pm 2m_i^2 I_2^i(0) \right) \pm m_i I_3^i(0) + I_1^i - \frac{1}{2G_3} = 0,
$$

(13)

$$
Z_{H,K} = \left( I_3^i(0) \pm 2m_i^2 I_2^i(0) \right)^{-1}.
$$

(14)

Note that eq. (13) is linear in $\Delta M$, and the $Z$–factors in eq. (14) are independent of $\Delta M$.

Let us also recall the numerical results following from our simple gradient expansion (13,14). While the values of the several mass differences $\Delta M_{H,K}^i$ are to be compared with experimental values, the $Z_{H,K}$–factors enter into the decay constants $f_{H,K}$ of heavy mesons defined by

$$
\langle 0|\bar{q}\gamma_\mu(1-\gamma_5)Q_v|H_v(0^-)\rangle = if_H M_H v_\mu,
$$

$$
\langle 0|\bar{q}\gamma_\mu(1-\gamma_5)Q_v|K_v(0^+)\rangle = -f_K M_K v_\mu,
$$

(15)

via the equation

$$
f_{H,K}\sqrt{M_{H,K}} = \frac{\sqrt{Z_{H,K}}}{G_3}.
$$

(16)

The light parameters $m_u,d = 300$ MeV, $m_s = 510$ MeV, $\Lambda = 1.25$ GeV, entering (13,14) are fixed from light meson properties, whereas the heavy–light quark coupling constant $G_3$ has to be adjusted from heavy meson physics.

As table 1 shows, it is impossible to find a consistent value of $G_3$ that fits both, the masses of heavy $H$ and $K$ mesons simultaneously. While a value of $G_3 = 8.7$ GeV$^{-2}$ gives realistic results for masses and decay constants of the members of the $H(0^−,1^-)$ multiplet \cite{4}, the masses and $Z$–factors for the parity conjugate mesons $K$ come out simply too large. This is in contrast to the light flavor sector where both vector and axial vector mesons are adequately reproduced. Obviously something must go wrong in the gradient expansion in the heavy sector. We therefore abandon the naive gradient expansion and re–investigate the exact expressions (9,10). This will lead us to a modified gradient expansion which overcomes the shortcomings of the naive gradient expansion in the heavy sector. It is important to note that such a modification of the gradient expansion should not be understood as a better approximation (recall that we even know the exact result) but rather as a heuristic method to extract physical information from quark loop calculations without including unphysical threshold effects. We hope that the following discussion will illuminate these ideas.

3 Modified gradient expansion

Let us now start to investigate the exact equations (9,10). Since we expect the mass shifts $\Delta M_K^i$ to be larger than $m_i^i$ we have to handle the imaginary part of the integral (12) above.
the threshold. For illustration let us drop the imaginary part and show the real part of eq. (9). In figure 2 we have plotted, for convenience, $G^{-1}$ as a function of $\Delta M_{H,K}$ resulting from the exact solution (9) and from the simple gradient expansion (13), respectively. (Heavy mesons with strangeness show a completely similar behaviour.) Note that already for values below the threshold $\Delta M = m$ the linear approximation seems problematic. Instead, $G^{-1}$ as a function of $\Delta M$ seems to start already with a significant curvature. Keeping more of the quadratic behaviour at $\Delta M = 0$ would immediately decrease $\Delta M_{K}$ for a given value of $G_{3}$ considerably, whereas $\Delta M_{H}$ would not such dramatically change.

For clarification let us compare the renormalization factors $Z_{i}^{H,K}$ in the exact formula (10) with the simple gradient expansion (14). As an example we plotted $(Z_{u}^{H})^{-1}$ in figure 3. Once again the simple gradient expansion cannot be viewed as a good low–energy approximation. This becomes even more apparent if we consider the analogous situation for the light sector. In figure 4 we plotted the real part of an exact calculation of the renormalization factor for the $\rho$ meson given by

$$Z_{\rho}^{-1}(p^{2}) = \frac{N_{c}}{16\pi^{2}} \int_{0}^{1} dx 4x(1-x)\Gamma(0, \frac{m^{2}-x(1-x)p^{2}}{\Lambda^{2}}) .$$

Here the usual gradient expansion is indeed justified as a definite interpolation from the low energy region to regions above the threshold.

Let us now try to understand analytically which terms are responsible for the deviation from the linear behaviour of eq. (9) (respectively from the constant behaviour of the $Z$–factors) in the low–energy region. The threshold presents itself in the combination $m^{2} - (x)\Delta M^{2}$ as one sees from the integral $I_{3}(v \cdot p) = \Delta M$ in eq. (12). For $\Delta M > m$ the incomplete gamma function and the square root with negative argument will produce imaginary parts. In the proper–time formalism this is the usual way how thresholds arise. Since the low–energy region is defined as $\Delta M \ll m$ we are on the safe side if we neglect $\Delta M^{2}$ compared to $m^{2}$ in the integral expressions, performing the replacement

$$I_{3}(\Delta M) \rightarrow \frac{N_{c}}{16\pi^{2}} \left\{ 2\Delta M \Gamma \left(0, \frac{m^{2}}{\Lambda^{2}}\right) + \sqrt{\pi}m \Gamma \left(-\frac{1}{2}, \frac{m^{2}}{\Lambda^{2}}\right) \right\}$$

$$I_{2}(\Delta M) \rightarrow I_{2}(0) .$$

Indeed such a strategy would reproduce a constant $Z$–factor for the $\rho$ meson. Furthermore, for heavy mesons the integral $I_{3}$ is multiplied by $(\Delta M \pm m)$, and this will generate a quadratic term (in $\Delta M$) in the self–energy. For light mesons such an additional pre–factor is missing. Note that our so defined modified gradient expansion has not to be viewed as an ordinary Taylor expansion of the self–energy. The crucial point of this investigation is that it would mean no improvement to calculate any higher order terms in the gradient expansion since we would get more and more unphysical effects from the threshold. Instead, our method represents a compromise for extracting relevant physics from the low energy region and neglecting unphysical imaginary parts and threshold effects leading to the emission of free quarks.

One could try to treat the thresholds in the quark loop integrals seriously and simply dropping imaginary parts. For the light (axial) vector mesons this would (accidently) make not much difference. However for heavy mesons such a strategy gets no numerical support and seems unphysical at all.

With this understanding we can dare to extrapolate from the region below threshold to mass shifts $\Delta M$ above threshold. Doing so we hope that chiral symmetry dominates the physics over a large region, at least so far that also the $(0^{+},1^{+})$ multiplet can be described by our modified
Figure 2: The behaviour of $1/G_3$ as a function of $\Delta M_{HH}/m^u$ (bold line) and of $\Delta M_{KK}/m^u$ (thin line). The relation (13) representing the lowest order gradient expansion is plotted in each case with the same line type and can be identified as a linear function originating from $\Delta M = 0$. 
Figure 3: The behaviour of the normalized $(Z_H^u)^{-1}$ as a function of $\Delta M_H^u/m^u$ (bold line). The threshold at $\Delta M/m = 1$ shows up as a singularity. In lowest order gradient expansion the $Z$–factor stays constant at the value at $\Delta M = 0$ (thin line).
Figure 4: The inverse normalized $Z$-factor $(Z_p)^{-1}$ as a function of $M_p^2/(4m^2)$ (bold line). This time the threshold at $M_p^2 = 4m^2$ shows up as a peak. The lowest order gradient expansion (thin line) is for a large region in good agreement with the exact solution.
gradient expansion. Equations (9), (10) are then to be replaced by

\[
\left( I_i^3(0) + 2I_i^2(0)\Delta M_{iH,K} \right) \left( \Delta M_{iH,K} \pm m_i \right) + I_i^1 \frac{1}{2G_3} = 0, \tag{20}
\]

\[
\left( Z_{iH,K}^i \right)^{-1} = I_i^3(0) + 2 \left( 2\Delta M_{iH,K} \pm m_i \right) I_i^2(0). \tag{21}
\]

Figure 5 shows the behaviour of \( G_3^{-1} \) in the modified gradient expansion \((20)\) to be compared with figure 2.

In table 2 we present the numerical results following from this modified gradient expansion. Note that heavy quark masses are extracted from \( m_{b,c} = M_{B,D} - \Delta M_H^{u} \). For the calculation of heavy meson decay constants we have used the experimental values of meson masses \cite{7} if available, or have made use of the relation \( M_{B_x} - M_{B} = M_{D_x} - M_{D} \) which is true in the heavy quark limit \( (x = \{s, 1, s_1\}) \). Indeed we observe a drastic improvement compared to the results in table 1. For a value \( G_3 \approx 7 \text{ GeV}^{-2} \) a simultaneous fit of nearly all heavy meson properties is possible. Only for the most massive \( (B_{s1}, D_{s1}) \) states, masses are still predicted slightly too high. With our analysis we are now able to predict the decay constant of several B–mesons (Note that \( f_D \) is assumed to get large \( 1/m_c \) corrections). Within a physical acceptable range of

\[
6.5 \text{ GeV}^{-2} < G_3 < 7.5 \text{ GeV}^{-2}
\]

we obtain

\[
160 \text{ GeV} < f_B < 180 \text{ GeV} \tag{23}
\]

\[
180 \text{ GeV} < f_{B_s} < 200 \text{ GeV} \tag{24}
\]

\[
150 \text{ GeV} < f_{B_1} < 160 \text{ GeV} \tag{25}
\]

\[
155 \text{ GeV} < f_{B_{s1}} < 175 \text{ GeV}. \tag{26}
\]

\footnote{6 This indicates that a simple local interaction is perhaps not sufficient to describe a heavy quark weakly bound via a p–wave with a not so light strange quark.}
Figure 5: The behaviour of $1/G_3$ as a function of $\Delta M_{H}/m^{u}$ (bold line) and of $\Delta M_{K}/m^{u}$ (thin line) using the modified gradient expansion (20) and the exact solution (9), respectively.
4 Conclusions

In this paper we have studied in detail the treatment of quark loop diagrams, arising from the calculation of the quark determinant in the recently proposed extended NJL model \cite{4}. Usually, the problems originating from the lack of confinement in quark models which leads to unphysical quark thresholds, are overcome by an expansion in external momenta. The lowest-order gradient expansion works successfully for composite light mesons. However, in the heavy meson sector there is phenomenological evidence that the gradient expansion has to be improved, which has been discussed in section 2.

We have shown that this fact is due to an essentially different analytical structure of the loop integrals under concern, coming from the special form of the heavy quark propagator \((v \cdot k + i\epsilon)^{-1}\).

While the renormalization factor \(Z_\rho\) for \(\rho\) mesons stays nearly constant below the threshold, this is not the case for the corresponding expressions \(Z_{H,K}\) of heavy mesons. We have convinced ourselves that this additional momentum dependence is not due to threshold effects and should be included. We stress again that our modified procedure should not be viewed as an ordinary Taylor expansion that would never be able to interpolate beyond the threshold.

Most importantly, with this improved gradient expansion heavy meson properties of both, the \((0^-,1^-)\) and the \((0^+,1^+)\) spin symmetry doublet can now be described simultaneously with a quark-coupling constant \(G_3 \approx 7\ \text{GeV}^{-2}\). With this value we are in a position to estimate several weak decay constants of B mesons.

In summary, it has been shown that the evaluation of quark loop diagrams with heavy quark propagators in unconfined models demands a modification of the gradient expansion that has been used in the light sector. A phenomenologically successful procedure seems to consist in neglecting unphysical two-quark threshold effects in a definite way and keeping all remaining momentum dependence. Phenomenologically, this leads to realistic results for heavy meson masses and decay constants within our extended NJL model with chiral and heavy quark symmetries.

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References

[1] N. Isgur and M. Wise, Phys. Lett. B232 (1989) 113; Phys. Lett. B237 (1990) 527; E. Eichten and B. Hill, Phys. Lett. B234 (1990) 511; B. Grinstein, Nucl. Phys. B339 (1990) 253; H. Georgi, Phys. Lett. B240 (1990) 447; T. Mannel, W. Roberts and Z. Ryzak, Nucl. Phys. B368 (1992) 204.

[2] W. A. Bardeen and C. T. Hill, Phys. Rev. D49 (1993) 409.

[3] M.A. Novak, M. Rho and I. Zahed, Phys. Rev. D48 (1993) 4370.

[4] D. Ebert, T. Feldmann, R. Friedrich and H. Reinhardt, Prepr. DESY 94–098, hep-ph/9406220 (submitted for publication).

[5] D. Ebert and M. K. Volkov, Yad. Fiz. 36 (1982) 1265; Z. Phys. C16 (1983) 205; M. K. Volkov, Ann. Phys. (N.Y.) 157 (1984) 282.

[6] D. Ebert and H. Reinhardt, Nucl. Phys. B271 (1986) 188; Phys. Lett. B173 (1986) 453.
[7] Particle Data Group, *Review of Particle Properties*, Phys. Rev. D 45 (1992).

[8] C. R. Allton et al., Nucl. Phys. B349 (1991) 598; C. Alexandrou et al., Phys. Lett. B256 (1991) 60; R. Sommer, DESY preprint 94–011 (1994).

[9] P. Colangelo, G. Nardulli, A. A. Ovchinnikov and N. Paver, Phys. Lett. B269 (1991) 201; S. Narison, Phys. Lett. B322 (1994) 247; CERN preprint TH. 7103/93; S. Narison and K. Zalewski, Phys. Lett. B320 (1994) 369.