Spin-dependent interaction in the deconfined phase of QCD

Yu.A.Simonov
State Research Center
Institute of Theoretical and Experimental Physics,
Moscow, 117218 Russia

November 7, 2018

Abstract

Spin-dependent deconfined interaction in the $Q\bar{Q}$ system is derived from the field correlators known from lattice and analytic calculations. As a result hyperfine splitting is found numerically for charmonium, bottomonium and strangeonium in the range $T_c \leq T \leq 2T_c$. Spin-orbit interaction due to magnetic correlators (the Thomas term) is able to produce numerous $Q\bar{Q}$ bound states with accumulation point at $M = m_Q + m_{\bar{Q}}$. Possible influence of these effects on the thermodynamics of quark-gluon plasma is discussed.

In honor of A.Di Giacomo on his seventieth birthday.

1 Introduction

The Spin-Dependent Interaction (SDI) in the $Q\bar{Q}$ system is being studied for the last three decades [1]-[12] for heavy quarkonia in the confinement phase. The first estimates of SDI [1, 5] essentially exploited perturbative expansion and nonperturbative SDI appeared in the form of the Thomas term [2]. In the general approach [3, 4] using $1/M_q$ expansion for heavy quarks, four spin-dependent static potentials $V_i(r), i = 1, 2, 3, 4$ are expressed through connected field correlators.
In this approach perturbative expansion of correlators reproduces perturbative part of $V_i(r)$ while the nonperturbative part can be obtained only from lattice simulations, done by several groups \[6\] \[7\] \[8\].

In the Field-Correlator Method (FCM) \[9\] (see \[10\] for a review) the SDI potentials have been obtained using $1/M_q$ expansion in \[11\] \[12\], in the form containing only Gaussian correlators $D(x)$, $D_1(x)$, and allowing for immediate check of Gromes relation \[4\]. The nonperturbative component of $V_i(r)$ is calculated directly through $D(x), D_1(x)$ and the Thomas term naturally appears as asymptotics at large distances. Extensive and thorough lattice calculations of correlators $D(x), D_1(x)$ done by the Pisa group \[13\] have made it possible to determine numerically all 4 potentials $V_i$.

The first comparison of resulting spin splittings of charmonium and bottomonium levels was done in \[14\] and has shown a good agreement of lattice-based potentials, with experiment. At the same time spin-orbit splittings were found to be sensitive to the behaviour of field correlators $D(x), D_1(x)$ at small $x$.

In \[15\] a detailed analysis was done of SD potentials, obtained from field correlators in \[11\], as compared to the definitions of Eichten-Feinberg-Gromes (EFG) \[3\] \[4\]. In the latter case SD potentials $V_i$ are not shown to satisfy Gromes relation \[4\], and analysis of \[15\] has demonstrated the difficulties arising in this respect for the lattice-defined $V_i$ \[6\] \[7\].

It was understood later \[16\], that the same set of SDI potentials $V_i(r)$ can be obtained without the $1/M_q$ expansion, using cluster expansion of gluonic fields for quarks of arbitrary mass $m_q$ and in this case instead of quark masses $m_q$, there appear in $V_i(r)$ the so-called einbein masses $\mu_q$, with effective average values $\langle \mu_q \rangle$ playing the role of constituent masses calculated through $m_q$ and string tension $\sigma$.

This approach works well in the confinement phase for meson \[16\] and baryon \[17\] states of light and heavy quarks, when SDI can be considered as corrections, see e.g. the analysis \[18\] for light scalar mesons, and \[19\] for heavy quarkonia.

The inclusion of chiral dynamics (where SDI is crucial) needs modification in the formalism, and instead of SDI potentials one considers in this case the effective Lagrangians, as in \[20\], and the huge mass splitting between $\rho$ and $\pi$ is obtained as the chiral nonlinear amplification of the standard hyperfine SDI.

So far so good for the zero temperature QCD. For $T > 0$ there appears a first change in SDI namely one must distinguish colorelectric and
colormagnetic correlators, contributing to $V_i(r)$, since instead of two Gaussian correlators $D(x), D_1(x)$ one has for $T > 0$ four independent correlators: $D^E(x), D^H(x), D^E_1(x), D^H_1(x)$ and it was agreed in [21, 22, 23] that only one of them vanishes for $T > T_c$. These correlators have been found numerically on the lattice [24], and the original expectation was confirmed, that the correlators do not change significantly for $T \ll M_d$, where $M_d$ is the dilaton mass of the order of $0^{++}$ glueball mass, $M_d \sim 1 \div 1.5$ GeV. Moreover it was shown in [24] that indeed only one of correlators, $D^E(x)$ vanishes exactly at $T = T_c$.

Already at this stage one can show that all four SDI potentials are expressed through magnetic correlators, and only one SDI term is due to the colorelectric correlator, contributing to spin-orbit force.

The situation becomes even more interesting for $T > T_c$, where the main colorelectric correlator $D^E(x)$, ensuring confinement, disappears but all other correlators stay intact.

In this case all four SDI potentials should change only slightly above $T_c$ and ensure the spin splitting of levels of the same type as it was below $T_c$. Of special interest is the Thomas term, which is doubled in magnitude (due to disappearance of colorelectric static force $\varepsilon(r)$) and is dominant at large distances. It is argued below in the paper, that the spin-orbit Thomas term can create infinite number of bound $Q\bar{Q}$ states with the accumulation point near $2M_Q$.

2 Spin-dependent interaction in the $q\bar{q}$ system

For the quark-antiquark Green’s function (of both light and heavy quarks) one can use the Fock-Feynman-Schwinger Representation (FSR) [25] with the kernel containing vacuum fields and quark spin operators in the form [11, 16]

$$\langle W_\sigma(x, y) \rangle = \langle \exp ig \int d\pi_{\mu\nu}(z) F_{\mu\nu}(z) \rangle =$$

$$= \exp \sum_{n=1}^{\infty} \frac{(ig)^n}{n!} \int d\pi(1) ... \int d\pi(n) \langle F(1)...F(n) \rangle,$$  

where

$$d\pi_{\mu\nu}(z) = ds_{\mu\nu}(z) - i\sigma_{\mu\nu}dt$$
Here $ds_{\mu\nu}$ is the surface element, while $d\tau$ is the proper time integration, which can be connected to usual time $t$ via $d\tau_i = \frac{dt_i}{2\mu_i}$ [16]. The spin-spin interaction can be obtained from [11, 2] keeping only quadratic (Gaussian) correlators $\langle FF \rangle$

$$\exp \left\{ -\frac{1}{2} \int_0^{s_1} d\tau_1 \int_0^{s_2} d\tau_2 g^2 \left\langle \left( \frac{\sigma^{(1)}B}{\sigma^{(1)}E} \right)_{z(\tau_1)} \left( \frac{\sigma^{(2)}B}{\sigma^{(2)}E} \right)_{z(\tau_2)} \right\rangle \right\},$$

and spin-orbit interaction arises in [11] from the products $\langle \sigma^{\mu\nu}ds^{\rho\lambda}F_{\rho\lambda} \rangle$ [11, 16].

It is clear, that the resulting SDI will be of matrix form $2\times2$ (not accounting for Pauli matrices). If one keeps only diagonal terms in $\sigma^{\mu\nu}F_{\mu\nu}$ (as the leading terms for large $\mu_i \approx M$) one can write for the SDI the representation of the Eichten-Feinberg form [26]

$$V_{SD}^{(diag)}(R) = \left( \frac{\sigma_1}{4\mu_1^2} - \frac{\sigma_2}{4\mu_2^2} \right) \left( \frac{1}{R} \frac{dz}{dR} + \frac{2dV_1(R)}{RdR} \right) +$$

$$+ \frac{\sigma_2 L_1 - \sigma_1 L_2}{2\mu_1\mu_2} \left( \frac{1}{R} \frac{dV_2}{dR} + \frac{3\sigma_2 R - \sigma_1 R^2}{12\mu_1\mu_2} \right).$$

At this point one should note that the term with $\frac{dz}{dR}$ in (4) was obtained from the diagonal part of the matrix $(m - \hat{D})\sigma_{\mu\nu}F_{\mu\nu}$, namely as product $i\sigma_k D_k \cdot \sigma_i E_i$; see [11, 16] for details of derivation, while all other potentials $V_i, i = 1, 2, 3, 4$ are proportional to correlators $\langle HH \rangle$. One can relate correlators of coloelectric and colormagnetic fields to $D^E, D_1^E, D^H, D_1^H$ as follow (see [11, 12])

$$g^2 \langle E_i(\vec{x}_1, \tau_1) \Phi E_j(\vec{x}_2, \tau_2) \Phi^+ \rangle = N_c \left[ \delta_{ij} \left( \frac{D^E(\lambda, \nu) + D_1^E + \vec{u}^2 \frac{\partial D_1^E}{\partial \vec{u}^2}}{u_i u_j \frac{\partial D_1^E}{\partial \vec{u}^2}} \right) u_i u_j \frac{\partial D_1^E}{\partial \vec{u}^2} \right]$$

$$g^2 \langle H_i(x) \Phi H_j(y) \Phi^+ \rangle = N_c \left[ \delta_{ij} \left( \frac{D^H(\lambda, \nu) + D_1^H + \vec{u}^2 \frac{\partial D_1^H}{\partial \vec{u}^2}}{u_i u_j \frac{\partial D_1^H}{\partial \vec{u}^2}} \right) - u_i u_j \frac{\partial D_1^H}{\partial \vec{u}^2} \right]$$

$$g^2 \langle E_i(x) \Phi H_j(y) \Phi^+ \rangle = -N_c e_{ijn} u_n \frac{\partial D_1^{EH}}{\partial \vec{u}^2}$$

As a result one obtains the following connection between SD potentials and correlators.
\[
\frac{1}{R} dV_1 = - \int_{-\infty}^{\infty} d\nu \int_0^R \frac{d\lambda}{R} \left( 1 - \frac{\lambda}{R} \right) D^H(\lambda, \nu) \tag{8}
\]

\[
\frac{1}{R} dV_2 = \int_{-\infty}^{\infty} d\nu \int_0^R \frac{\lambda d\lambda}{R^2} \left[ D^H(\lambda, \nu) + D^H_1(\lambda, \nu) + \lambda^2 \frac{\partial D^H_1}{\partial \lambda^2} \right] \tag{9}
\]

\[
V_3 = - \int_{-\infty}^{\infty} d\nu R^2 \frac{\partial D^H_1(R, \nu)}{\partial R^2} \tag{10}
\]

\[
V_4 = \int_{-\infty}^{\infty} d\nu \left( 3D^H(R, \nu) + 3D^H_1(R, \nu) + 2R^2 \frac{\partial D^H_1}{\partial R^2} \right) \tag{11}
\]

\[
\frac{1}{R} d\varepsilon(R) = \int_{-\infty}^{\infty} d\nu \int_0^R \frac{d\lambda}{R} \left[ D^E(\lambda, \nu) + D^E_1(\lambda, \nu) + (\lambda^2 + \nu^2) \frac{\partial D^E_1}{\partial \nu^2} \right] \tag{12}
\]

One can check, that the Gromes relation acquires the form

\[
V'_1(R) + \varepsilon'(R) = V'_2(R) = \int_{-\infty}^{\infty} d\nu \left[ \int_0^R d\lambda (D^E(\lambda, \nu) - D^H(\lambda, \nu)) + \frac{1}{2} R (D^E_1(R) - D^H_1(R)) \right]. \tag{13}
\]

For \( T = 0 \), when \( D^E = D^H, D^E_1 = D^H_1 \), the Gromes relations are satisfied identically, however for \( T > 0 \) electric and magnetic correlators are certainly different and Gromes relation is violated, as one could tell beforehand, since for \( T > 0 \) the Euclidean \( O(4) \) invariance is violated.

To conclude this section, one comment on nondiagonal terms in (3), which contribute to the total Hamiltonian \( \hat{H} \) as

\[
\hat{H} = H_0(\mu_1, \mu_2) + V_{SD}^{(diag)} \hat{1}_1 \hat{1}_2 + V_{SD}^{(nond)}(\gamma_5)_{1}(\gamma_5)_{2} + \ldots \tag{14}
\]

where dots imply terms proportional to \( \hat{1}_1(\gamma_5)_{2} \) and \( (\gamma_5)_{1} \hat{1}_2 \). It is clear that \( V_{SD}^{(nond)} \) has the structure similar to (4) with replacement \( V_i \to V_i^{(nond)} \) and contains terms proportional to electric correlators Eq.(5). E.g. in (26) the term \( V_i^{(nond)} \) was found to be

\[
V_i^{(nond)}(R) = \int_{-\infty}^{\infty} d\nu \left( 3D^E + 3D^E_1 + (3\nu^2 + R^2) \frac{\partial D^E_1}{\partial R^2} \right) \tag{15}
\]

For heavy quarks (and for light quarks in the states, where SDI is repulsive) the extremal values \( \mu_i^{(0)} \) can be found from the minimum of the spinless
Hamiltonian $H_0(\mu_1, \mu_2)$, and in this case SDI gives spin corrections, which are not large even for light quarks and are in good agreement with experiment – see [18, 19] for heavy and light quarkonia respectively.

For light quarks in the states with attractive spin-spin interaction the $q\bar{q}$ Hamiltonian should be taken with the full matrix structure as in (14), and the stationary values of the quark "constituent" masses $\mu_i^{(0)}$ should be found from the minimum of eigenvalues of $\hat{H}$, see [26] for more details. In what follows we shall consider only diagonal part of SDI, which is essential for heavy quarkonia.

3 Spin interactions in the deconfinement phase

In the confined phase SDI in terms of field correlators was studied in [11, 12, 15–19]. For $T > T_c$ the correlator $D^E(x)$ vanishes, as it was argued in [21, 22, 23] and confirmed in lattice calculations [24]. This fact leads to a serious change both in SDI as well as in the spin-independent part $H_0$ of the total Hamiltonian $\hat{H}$. The latter change was studied in detail in recent papers [27, 28], where it was shown that in the total static potential $V(R) = V_D(R) + V_1(R)$, generated by $D^E(x)$ and $D^E_1(x)$ respectively, only the term $V_1(R)$ survives for $T > T_c$. This term in contrast to the confining potential $V_D(R) \sim \sigma R, \ R \to \infty$, saturates at large $R$ and can support bound states of $c\bar{c}$ and $b\bar{b}$ in some interval of temperatures $T_c \leq T \leq T_d$, while bound states dissociate at $T \geq T_d$. ($T_d \sim 2T_c$ for $D^E_1$ calculated analytically in [29], and $T_d \sim 1.12T_c$ for $D^E$ from lattice calculations done in [24], see [28] for details). All this dynamics is due to colorelectric correlator $D^E$ and colormagnetic fields were not taken into account. We now consider the role of SDI, which is mostly due to colormagnetic fields.

We start with small distances and assume, according to [27, 28, 29] the following form of $D^E_1(D^H_1, D^H)\,$

$$D^{E,H}_1 = D^{E,H}_1(\text{pert}) + D^{E,H}_1(\text{nonpert}), \quad D^H = D^H_{\text{nonpert}} + D^H_{\text{pert}},$$

(16)

where

$$D^{E,H}_1(\text{pert}) = \frac{16\alpha_s}{3\pi x^4} e^{-\lambda_{E,H}x} + O(\alpha_s^2); \quad D^{E,H}_1(\text{nonpert}, M_{E,H} x \gg 1) = C_{E,H} \frac{e^{-M_{E,H} x}}{x}$$

(17)

and

$$D^H_{\text{nonpert}}(x M_0 \gg 1) = d_H e^{-M_0 x}, \quad D^H_{\text{pert}} = \frac{g^4}{4\pi^4 x^4} + O(\alpha_s^3)$$

(18)
Here \( C_{E,H}, d_H \) are known constants depending on \( \sigma_H \) - the spacial string tension, \( \sigma_H = \frac{1}{2} \int d^2 \lambda D_{nonpert}^H(\lambda, \nu) d^2 x \).

One can easily check in \( [8,12] \), that the small distance behaviour of SDI potentials \( V_i(R) \) does not change much, since it is defined by the perturbative parts of correlators which are little modified when temperature grows above \( T_c \). The main difference comes at large distances, where according to Eqs. \( [8,12] \) one finds in the limit of large \( R \),

\[
V_1'(R) = -\sigma_H, \quad V_2'(R) = \frac{\gamma_H}{R}, \quad \gamma_H = \int_{-\infty}^{\infty} d\nu \int_{0}^{\infty} \lambda d\lambda D_{nonpert}^H(\lambda, \nu)
\]

\[
\varepsilon'(R) = \frac{1}{2} \int_{-\infty}^{\infty} d\nu RD_1^E(R) \to 0, \quad V_3(R) \to 0, \quad V_4(R) \to 0
\]

where in \( [20] \) all three potentials are exponentially small at large \( R \). Therefore the spin-orbit part of \( V_{SD}^{(diag)} \) has the form (for equal quark pole masses)

\[
V_{SD}^{(diag)}(R \to \infty) = -\frac{\text{SL} \sigma_H}{\mu^2 R} + \frac{\text{SL} \gamma_H}{\mu^2 R^2} + O(e^{-MR}).
\]

The first term on the r.h.s. of \( [21] \) was quoted (without derivation) in \( [21] \), where it was suggested that being dominant at large \( R \) (where \( V_1(R) \) exponentially approaches to the constant limit \( V_1(\infty) \)) this term by itself can support bound states of heavy (and, possibly, light) quarkonia. Indeed, considering the first term in \( [21] \) as a Coulomb-like potential one arrives for heavy quarkonia (where \( \mu_0 \approx m \) ) to the mass spectrum of bound states for \( \text{SL} > 0, J = L + 1 \),

\[
M_n = 2m - \frac{\sigma_H^2}{4m^2} \frac{L^2}{(L + n + 1)^2}, \quad n = 0, 1, 2, ...
\]

We conclude this section with discussion of hyperfine interaction in the deconfined phase. As it is seen in \( [11] \), \( V_4(R) \) depends only on \( D^H, D_1^H \) and does not change across \( T_c \) according to the lattice data \( [24] \). Keeping only (the dominant) perturbative part of \( V_4 \), one obtains

\[
V_{hf} = \frac{8\pi \alpha_s}{9\mu_1 \mu_2} \delta^{(3)}(R).
\]

For the hyperfine energy shifts this gives

\[
\Delta E_{hf} = \frac{4\alpha_s}{9(\mu_1 + \mu_2)} \begin{pmatrix} -3, & S = 0 \\ +1, & S = 1 \end{pmatrix}.
\]
To calculate \( \langle V'(R) \rangle \), one can use the static potential, computed in \[27, 28\] both analytically and on the lattice. A rough estimate can be found from comparison of free static energies found on the lattice (see e.g. \[30\] and refs. therein) both below and above \( T_c \) (up to \( T \approx 1.3T_c \)), which shows a close similarity of \( V'(R) \) in both temperature domains. Hence one expects for \( T_c \leq T \leq 1.3T_c \) the mass gap between \( J/\psi \) and \( \eta)c \) of the same order as for \( T = 0 \), i.e. \( \delta E \sim 0.1 \) GeV, which roughly agrees with lattice MEM computations \[30\].

The situation with hyperfine spin splittings in the light \( q\bar{q} \) system is however different and cannot be treated by the methods given above, since the restoration of chiral symmetry at \( T \geq T_c \) needs the effective Lagrangian technic \[20\] mentioned above. The physical reason for that lies in the fact, that for light quarks \( \Delta E_{hf} \) Eq. \[24\], becomes dominant (for \( \mu_1 = \mu_2 \approx m \)) and cannot be treated as a perturbation, and one needs to solve nonlinear equations for the effective mass operator, given in \[20\]. This point will be treated elsewhere.

In summary, we have derived the spin-dependent potentials for the \( q\bar{q} \) system for any temperature \( T \) valid in the situation, when spin splittings can be considered perturbatively.

For \( 1.3T_c \geq T \geq T_c \) spin splittings of charmonium and bottomonium are shown to change little compared to zero temperature case. The color magnetic confinement produces the Thomas spin-orbit term, which dominates at large distances, in absence of colorelectric confinement at \( T \geq T_c \), and can possibly support a sequence of bound states. These features demonstrate the importance of strong interaction in the quark-gluon plasma, which was advocated in \[21-24\] and supported by explicit calculations in \[27, 28\] in agreement with lattice data \[30\].

This work is supported by the Federal Program of the Russian Ministry of industry, Science and Technology No.40.052.1.1.1112, and by the grant for scientific schools NS-1774. 2003.

References

[1] A.De Rujula, H.Georgi and S.L.Glashow, Phys. Rev. D12, 147 (1975); T.de Grand, P.L.Jaffe, K.Johnson and J.Kiskis, Phys. Rev. D12, 2060 (1975).
[2] T.Appelquist, R.M.Barnett and K.D.Lane, Ann. Rev. Nucl. Part. Sci. 28, 387 (1978);
W.Buchmueller, Phys. Lett. 112, 479 (1982).

[3] E.Eichten, F.Feinberg, Phys. Rev. D23, 2724 (1981).

[4] D.Gromes, Z.Phys. C26, 401 (1984).

[5] W.Buchmüller, Y.J.Ng and S.-H.H.Tye, Phys. Rev. D24, 3003 (1981);
J.Pantaleone, Y. J.Ng and S.-H.H.Tye, Phys. Rev. D33, 777 (1986).

[6] M.Campostrini, Nucl. Phys. B256, 717 (1985);
C.Michael and P.E.L.Rakow, Nucl. Phys. B256, 640 (1985);
P.de Forcrand and J.D.Stack, Phys. Rev. Lett. 55, 1254 (1985);
M.Campostrini, K.Moriarty and C.Rebbi, Phys. Rev. Lett. 57, 44 (1986);
C.Michael, Phys. Rev. Lett. 56, 1219 (1986);
I.J.Ford, J.Phys. G 15, 1571 (1989);
A.Huntley and C.Michael, Nucl. Phys. B270, 123 (1986).

[7] K.D.Born et al. Phys.Lett. B329, 332 (1994);
G.S.Bali, K.Schilling and a.Wachter, Phys. Rev. D55, 5309 (1997), ibid. D56, 2566 (1997).

[8] M.Koma, Y.Koma, H.Wittig, hep-lat/0510059.

[9] H.G.Dosch, Phys. Lett. B 190, 177 (1987) ;
H.G.Dosch and Yu.A.Simonov, Phys. Lett. B 205, 339 (1988) ;
Yu.A.Simonov, Nucl. Phys. B 307, 512 (1988).

[10] A.Di Giacomo, H.G.Dosch, V.I.Shevchenko, Yu.A.Simonov, Phys. Rept. 372, 319 (2002); hep-ph/0007223.

[11] Yu.A.Simonov, Nucl. Phys. B324, 67 (1989).

[12] M.Schiestl, H.G.Dosch, Phys. Lett. B209, 85 (1988).

[13] A.Di Giacomo and H.Panagopoulos, Phys. Lett. B285, 133 (1992) ;
M.D’Elia, A.Di Giacomo, and E.Meggiolaro, Phys. Lett. B408, 315 (1997) ;
A.Di Giacomo, E.Meggiolaro and H.Panagopoulos, Nucl. Phys. B483,
371 (1997)
M.D’Elia, A.Di Giacomo, and E.Meggiolaro, Phys. Rev. D67, 114504 (2003);
G.S.Bali, N.Brambilla, A.Vairo, Phys. Lett. B 42, 265 (1998);
E.Meggiolaro, Phys. Lett. B451, 414 (1999).

[14] A.M.Badalian, V.P.Yurov, Yad. Fiz. 56, 239 (1993).

[15] A.M.Badalian, Yu. A.Simonov, Phys. At. Nucl. 59, 2164 (1996).

[16] Yu.A.Simonov, QCD and Topics in Hadron Physics, Lectures at the
XVII International School of Physics, Lisbon, 29 September -4 October,1999. [hep-ph/9911237]
Yu.A.Simonov, Spin-interactions of light quarks, preprint ITEP 97-89
(unpublished).

[17] Yu.A.Simonov, Phys. Rev. D65, 116004 (2002).

[18] A.M.Badalian, Phys. At. Nucl. 66, 1342 (2003);
A.M.Badalian, B.L.G.bakker, Phys. Rev. D67, 071901 (2003).

[19] A.M.Badalian, V.L.Morgunov, B.L.G.Bakker, Phys. At. Nucl. 63, 1635
(2000).

[20] Yu.A.Simonov, Phys. Rev. D65, 094018 (2002);
S.M.Fedorov, Yu.A.Simonov, JETP Lett. 78, 57 (2003).

[21] Yu.A.Simonov, JETP Lett. 54, 249 (1991).

[22] Yu.A.Simonov, JETP Lett. 55, 605 (1992).

[23] Yu.A.Simonov, Phys. At. Nucl. 58, 309 (1995); [hep-ph/9311216]
N.O.Agasian, JETP Lett. 57, 208 (1993).

[24] M. D’Elia, A. Di Giacomo and E. Meggiolaro, Phys. Rev. D 67, 114504
(2003);
A. Di Giacomo, E. Meggiolaro and H. Panagopoulos, Nucl. Phys. B 483,
371 (1997);
A. Di Giacomo, E. Meggiolaro and H. Panagopoulos, [hep-lat/9603018]

[25] Yu.A.Simonov, J.A. Tjon, Ann. Phys. 300, 54 (2002); [hep-ph/0205165]
Yu.A.Simonov and J.A.Tjon, Ann Phys. 228, 1 (1993).
[26] Yu.A.Simonov, Phys. At. Nucl. 68, 709 (2005).

[27] Yu.A.Simonov, Phys. Lett. B619, 293 (2005), hep-ph/0502078.

[28] A.Di Giacomo, E.Meggiolaro, Yu.A.Simonov, A.I.Veselov, hep-ph/0512125

[29] Yu.A.Simonov, Phys. At. Nucl. 69, N3, (2006), hep-ph/0501182

[30] P.Petreczky, hep-lat/0409139