Oscillatory behavior of the in-medium interparticle potential in hot gauge system with scalar bound states

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We investigate the in-medium interparticle potential of hot gauge system with bound states by employing the QED and scalar QED coupling. At finite temperature an oscillatory behavior of the potential has been found as well as its variation in terms of different free parameters. We expect the competition among the parameters will lead to an appropriate interparticle potential which could be extended to discuss the fluid properties of QGP with scalar bound states.

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I. INTRODUCTION

When concerning a hot/dense plasma, Debye screening picture is widely accepted where the in-medium screening potential is exponentially damping with distance. The damping rate, namely the so-called Debye mass, is the reciprocal of the screening length which scales the effective interactive radius of in-medium particles. At high temperature, the mean free path of particle is much larger than the screening length, therefore such kind of system can be treated as weakly coupled gas composed of quasi-particles with effective radius. Nevertheless this short-distance dominant Debye screening is not the whole story. Due to the dissimilar analytic structures of the effective boson propagator, the screening potential may behave distinctly at different situations and approximations, in which large distance oscillations become equally significant.

To make it more pellucid, we follow notations in Ref.[1] and defined the in-medium interparticle potential $V(r)$ in linear response framework as

$$V(r) = \frac{Q_1 Q_2}{4\pi^2 r} \Im \int_{-\infty}^{\infty} dq \frac{q e^{iqr}}{q^2 + F(0, q)},$$

where $r$ is the distance between two particles, $Q_1$ and $Q_2$ are their charges, $Q = (\omega, \mathbf{q})$ is the four-momentum transfer with $q = |\mathbf{q}|$, and $F$ is the longitudinal boson self-energy defined by $F = (Q^2/q^2)\Pi_{00}$. Obviously, the behavior of $V(r)$ is totally determined by the analytic structures of the denominator of the integrand on the complex plane. The denominator and the so-called dielectric function are almost the same except for a difference of $Q^2$. So we can simply say that the form of dielectric function determines the interactive potential.

Take the Friedel oscillations and Debye screening as examples. In the extremely high density but zero temperature environment, the explicit expressions of $F$ for nuclear matter[2, 4] and gauge plasma[2] have been obtained by ignoring the fermion masses. Due to the existence of sharp fermi surface, the dielectric function contains branch cuts, which dominate the contribution of contour integration at large distance and lead the interparticle potential oscillate considerably instead of monotonic damping. This phenomenon induced by the sharp fermi surface is named as Friedel oscillations[1, 2, 3, 5]. However at non-zero temperature the Friedel oscillations will be strongly suppressed because of the smeared Fermi surface at finite temperature. Fortunately, although the branch cut contribution is suppressed at finite temperature, the pole contribution stands out, resulting in the so-called Yukawa oscillations[1, 6] which dominate in a wide range of distance. On the contrary, at high temperature but zero chemical potential, there is only pole but no branch cut contribution. At hard thermal loop(HTL) approximation, the poles of the integrand are on the imaginary axis[6] which grants Debye screening. While in a not very high temperature environment, the HTL approximation might not be eligible and thus one may try to give up this approximation and employ a complete one-loop self-energy. However, the analytic structures of the self-energy are much more complicated in one-loop order. Actually, the position of the pole contains both imaginary and real parts on the complex plane which can only be identified numerically. The emergence of the real part of the pole means the interparticle potential is not monotonically

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damping but oscillating instead. For example, if there is a pole in the integrand of Eq. (1), whose imaginary and real parts are defined as \( q_i \) and \( q_r \) respectively, the potential can be obtained by performing the contour integral in the upper plane and recast as

\[
V(r) = \frac{Q_1 Q_2}{\pi(a^2 + b^2)} \frac{e^{-q_r r}}{r} \left[ a \cos(q_r r) + b \sin(q_r r) \right],
\]

where \( a \) and \( b \) are defined as the real and imaginary parts of the residue,

\[
\frac{(q^2 + F(q))'}{q} \bigg|_{q=q_r+iq_i} = a + ib,
\]

with the prime denoting for \( \partial/\partial q \). Considering the real part of \( F \) is an even function, we can only concern the residue in the first quadrant and double it in the final result. From mathematical points of view, when the pole is on the imaginary axis, the oscillating factor in the blanket of Eq. (2) vanishes as a result of \( q_r = 0 \), which leads to Debye screening potential. In other words, if the pole deviates from the imaginary axis, there exit oscillations.

Besides the mathematical description of the oscillations, we have an another potential motivation on the discussion of the oscillatory potential. As we all know the quark-gluon plasma (QGP) produced at relativistic heavy ion collider (RHIC) is almost likely to be perfect fluid at temperatures in the order of several times of critical temperature \( T \sim (1-3)T_c \). The fascinating low viscous mechanism has attracted much attention but still challenges both experimentalists and theoreticians. One possible way to understand it is to calculate the shear viscous coefficient. In the weakly-coupled regime many publications such as Refs. [7, 8, 9, 10, 11, 12, 13, 14, 15, 16, 17, 18, 19, 20, 21, 22, 23, 24, 25], no matter from kinetics or Kubo formula, are overestimating the shear viscosity extracted from experiments or lattice results [26, 27, 28]. The failure prompts to examine the picture of weak coupling and short-distance screening potential, and consider the strong coupling and/or long-distance correlation. Apart from the attempts in weak coupling, some scientists proposed strongly-coupled plasma and the liquid state theory [23, 30] to explain the perfect fluid QGP. One should keep in mind that the necessary condition to form a liquid state is the interparticle potential is strong enough in coupling as well as long enough in distance [31, 32]. Obviously Debye potential is not a good candidate because of its short distance domination character. Therefore it makes sense to study the oscillatory potential with respect to all kinds of dynamical and thermodynamical parameters so as to make sure if it is possible to produce a liquid state in such a hot environment.

Shuryak and Zahed [33] raised the idea of colored loosely binary bound states to explain the small viscosity at the near critical temperatures. In their picture, quark and anti-quark are bounded together, leaving a weak Coulomb potential as a remnant for interacting with the colored particles, such as quarks, gluons or other bound states in QGP. They argued the possible existing low-lying resonances led to large “unitarity limit” scattering cross section and naturally were expected to play an important role in the transport coefficients. In this paper, we adopt this picture of relativistic plasma with bound states, attempting to study their effect on the in-medium interparticle potential by using QED and scalar QED (sQED) as toy models. Since we do not intend to describe the real electromagnetic system but just quote their interaction forms to see the qualitative properties of the in-medium interparticle potential, so the temperature, the coupling strength and the masses of fermion and scalar bound state, are free parameters in the following calculation.

II. QED AND SQED SELF-ENERGY

In Eq. (1), one can see all the complication comes from the longitudinal self-energy of gauge boson. In one-loop order the photon self-energy is diagrammatically denoted by Fig. 1(a) in pure QED plasma and modified by adding Figs. 1(b) and (c) according to sQED Feynman rules when considering the existence of scalar bound states. The explicit expression of \( F \) for Fig. 1(a) is composed of vacuum and matter contributions, that are \( F_{\text{vac}} \) and \( F_{\text{mats}} \) respectively.
The vacuum contribution can be obtained by dimensional regulation and \(\overline{MS}\) scheme. The result can be found in many textbooks on Quantum Field Theory\cite{34},

\[
F^a_{\text{vac}}(0, q) = -\frac{2\alpha}{\pi} q^2 \int_0^1 dx (1-x) \log \left( \frac{m^2}{m^2 + x(1-x)q^2} \right),
\]

where \(\alpha = \frac{1}{137}\) is the fine structure constant of QED, and \(m\) is the mass of fermion. The matter contribution can be obtained either in real time or imaginary time formalism which reads\cite{35}

\[
F^a_{\text{mat}}(0, q) = -\frac{4\alpha}{\pi} \int_0^\infty dp \frac{p^2}{\omega} \left[ 4\omega^2 - q^2 \frac{\log \left( \frac{q - 2p}{q + 2p} \right) - 1}{4p} \right] n_f(\omega),
\]

where \(\omega\) is the on-shell energy of fermion constrained to \(\omega = \sqrt{p^2 + m^2}\) and \(n_f(\omega) = (e^{\beta\omega} + 1)^{-1}\) is the Fermi-Dirac distribution function with \(\beta = 1/T\).

The existing scalar bound states will "supplement" the polarization of the gauge field and consequently remodel the interaction among particles. We involve this effect by employing scalar QED. The polarization and tadpole diagrams presented in Figs.1(b) and (c) will be added to (a). In the real time formalism \(F\) can be written as\cite{36, 37}

\[
F^b(0, q) = -2i\pi\alpha_s \int \frac{d^4 P}{(2\pi)^4} (2p_0 - k_0)^2 [\Delta_F(P - K) \Delta_R(P) + \Delta_A(P - K) \Delta_F(P)],
\]

\[
F^c(0, q) = 4i\pi\alpha_s \int \frac{d^4 P}{(2\pi)^4} [\Delta_R(P) + \Delta_A(P) + \Delta_F(P)],
\]

where \(\alpha_s\) is the effective strength for the coupling with scalar bound states, capital letters denote for four-momentum, \(\Delta_F, \Delta_R\) and \(\Delta_A\) are the propagators in Keldysh representation\cite{38, 39} which read,

\[
\Delta_{R,A}(P) = \frac{1}{P^2 - m_s^2 \pm isgn(p_0)\varepsilon},
\]

\[
\Delta_F(P) = -2\pi i [1 + n_b(|p_0|) \delta(P^2 - m_s^2)],
\]

where \(m_s\) is the scalar mass, \(\varepsilon\) is the infinitesimal, \(sgn\) denotes for the sign function and \(n_b(|p_0|) = (e^{\beta|p_0|} - 1)^{-1}\) is the Bose-Einstein distribution function.

Separating the vacuum contribution from Eqs.\(6\) and \(7\), one can obtain

\[
F^{b+\varepsilon}_{\text{vac}}(0, q) = -\frac{\alpha_s}{4\pi} \int_0^1 dx (1-x)^2 \log \left( \frac{m_s^2}{m_s^2 + x(1-x)q^2} \right),
\]

\[
F^{b+\varepsilon}_{\text{mat}}(0, q) = -\frac{4\alpha_s}{\pi} \int_0^\infty dp \frac{p^2}{\omega^2} \left[ \frac{\omega^2}{p} \log \left( \frac{q - 2p}{q + 2p} \right) - 1 \right] n_b(\omega_s),
\]

where the vacuum contribution is manipulated in the same renormalization scheme as Eq.\(1\), where \(\omega_s = \sqrt{p^2 + m_s^2}\).

**III. NUMERICAL RESULTS**

To see the oscillatory behavior of the interparticle potential, we begin with the pure QED coupling case, i.e., we consider only Fig.1(a) in \(F\). In the following numerical calculations, only attractive interaction between two fermions with opposite sign is studied, so that \(Q_1 Q_2 = -4\pi\alpha\).

Inserting the sum of Eqs.\(1\) and \(4\) into Eq.\(1\), the pole of the integrand in the first quadrant can be identified numerically by replacing all \(q\) as \(q_r + iq_i\) in the denominator and then solving the complex equation. After that, one can calculate the undetermined parameters \(a\) and \(b\) through Eq.\(3\) and consequently figure out the potential according to Eq.\(2\). In Fig.2 we demonstrated the real part \(q_r\) (solid) and imaginary part \(q_i\) (dashed) of the pole as functions of temperature and the fermion mass. In Fig.2(a), the mass has been fixed at 0.5MeV and in Fig.2(b) the temperature has been fixed at 0.2GeV. According to Eq.\(2\), the fact \(q_r\) and \(q_i\) increase with temperature indicates...
FIG. 2: Position of poles in terms of temperature and fermion mass in pure QED coupling, where the solid lines are for $q_r$ and the dashed lines are for $q_i$. The curves in (a) are plotted with fixed fermion mass $m = 0.5\text{MeV}$ and those in (b) are plotted with fixed temperature $T = 0.2\text{GeV}$.

FIG. 3: Pure QED interparticle potential oscillates in distance. The curves in (a) are for fixed mass $m = 0.5\text{MeV}$ with the temperatures are $0.8\text{GeV}(\text{dotted})$, $0.4\text{GeV}(\text{dashed})$ and $0.2\text{GeV}(\text{solid})$ respectively. The curves in (b) are for fixed temperature $T = 0.2\text{GeV}$ with the effective fermion masses are $0.5\text{MeV}(\text{solid})$, $0.8\text{MeV}(\text{dashed})$ and $1.6\text{MeV}(\text{dotted})$ respectively.

the oscillation damps faster and more rapidly with increasing temperature. As to the fermion mass, there exist an opposite tendency, i.e., fast damping and rapid oscillations will be achieved by decreasing the mass.

To make this analysis more visible, we choose three points on each curve in Fig.2 to plot the oscillatory potential in Fig.3 where Fig.3(a) is fixed at $m = 0.5\text{MeV}$ with the temperatures are $0.8\text{GeV}(\text{dotted})$, $0.4\text{GeV}(\text{dashed})$ and $0.2\text{GeV}(\text{solid})$ respectively. Fig.3(b) is fixed at $T = 0.2\text{GeV}$ with the effective fermion masses are $0.5\text{MeV}(\text{solid})$, $0.8\text{MeV}(\text{dashed})$ and $1.6\text{MeV}(\text{dotted})$ respectively. Note that the solid curve in the two plots share the same set of parameters, and thus can be chosen as a reference.

In Fig.3 one can see clearly the interparticle potential oscillates in distance, where the peaks and the oscillating damping tail vary with temperature and mass. The first peak, which is the most obvious one on the plot and can be regarded as the representation of all the peaks, decreases in amplitude and becomes broad in width with increasing temperature, which means the effective interacting distance becomes long though weak at relatively low temperature. Contrary to the temperature effect, the decreasing fermion mass enhances and narrows the peak as shown in Fig.3(b).

Now we are in the position to present the effect of bound states. To involve this effect, the contributions from the last two diagrams in Fig.1 should be added to $F$, i.e.,

$$F = F_{\text{vac}}^a + F_{\text{matt}}^a + F_{\text{vac}}^{b+c} + F_{\text{matt}}^{b+c}.$$  \hspace{1cm} (11)

Numerically figuring out the zero point of $q^2 + F$ in Eq.(11), one can obtain the oscillatory potential with respect to different group of parameters. In Fig.4 we exhibit $q_r(\text{solid})$ and $q_i(\text{dashed})$ in terms of $T$, $\alpha_s$, $m$ and $m_s$. The parameters we used in the calculation are presented below each plot. As a general impression, one can find out two tendencies: one is ascending with increasing parameters, like the first row in Fig.4. The other is descending with increasing parameters, like the second row. Based on the previous experience on pure QED coupling, one could expect the first tendency would enhance but narrow the peaks of interparticle potential and shift them to the left, which means strong but short distance interaction. While the second tendency would suppress but broaden the peaks and shift them to the right, which suggests the weak but long distance interaction. We show these tendencies more clearly
\[ \alpha = \frac{1}{137}, \alpha_s = 0.02, m_s = 0.4\text{GeV}, m = 10\text{MeV} \]

\[ \alpha = \frac{1}{137}, T = 0.5\text{GeV}, m_s = 0.4\text{GeV}, m = 10\text{MeV} \]

\[ \alpha = \frac{1}{137}, \alpha_s = 0.02, m_s = 1.5\text{GeV}, m = 10\text{MeV} \]

\[ \alpha = \frac{1}{137}, \alpha_s = 0.02, T = 0.5\text{GeV}, m = 10\text{MeV} \]

FIG. 4: Position of poles in terms of \( T, \alpha_s, m \) and \( m_s \), where the solid lines are for \( q_r \) and the dashed lines for \( q_i \) as in Fig.2. The parameters we used in the calculation are presented below each plot.

in Fig.3 by directly exhibiting the potential in terms of these parameters where we set: \( \alpha = \frac{1}{137}, m = 10\text{MeV} \) and \( T = 0.5\text{GeV} \).

One can observe that Fig.3 is consistent with what we have argued by just analyzing the evolutional tendencies of the pole position with respect to various parameters. In each plot of Fig.3 the two curves denote for the case of \( m_s = 1.5\text{GeV} \) (solid) and \( m_s = 0.3\text{GeV} \) (dashed) respectively, which shows the decreasing of \( m_s \) or increasing of \( \alpha_s \) enhances but narrows the peaks on the potential and shift them to the left, indicating the strong but short distance interaction. Besides, it is worthwhile to compare the differences between the two curves in the same plot from the upper and lower ones which are mapped in different \( \alpha_s \). The comparison shows the stronger the coupling \( \alpha_s \) is, the more the scalar mass can affect the shape of the potential, even the other parameters are the same.

Finally let us compare the pure QED with the case involving scalar bound states. In Fig.6 we demonstrated the potential at \( \alpha_s = 0 \) (pure QED) and \( \alpha_s = 0.05 \) in upper plot with \( T = 0.3\text{GeV} \) and arrayed it with the lower one at \( T = 0.6\text{GeV} \). The masses of the fermion and bound state are 10MeV and 0.4 GeV respectively. There is no surprise that the pure QED potential is longer in the interactive range but weaker in amplitude than that involves the coupling with scalar bound state since it is the limit of the small coupling strength.

IV. SUMMARY AND DISCUSSIONS

We investigated the in-medium interparticle potential of gauge plasma with bound states by employing QED and sQED coupling, where complete one-loop boson self-energy has been involved which is equivalent to resum all one-loop diagrams. Different from Debye screening via HTL approximation, an oscillatory behavior of the interparticle potential with exponential damping amplitude has been demonstrated which is referred to the so-called Yukawa oscillation. The four parameters, \( \alpha_s, T, m, \) and \( m_s \) in our model are free to change, since we do not intend to study the electromagnetic system but to employ it as a toy coupling form. In the numerical computation we fixed \( \alpha \) to the fine structure constant by considering the matter particle as point-like. While \( \alpha_s \) was treated as an effective parameter because of the inner structure of the scalar bound state. Evidently, \( \alpha_s = 0 \) recovers the pure QED coupling. In the
FIG. 5: In-medium interparticle potential varies with the scalar mass at different $\alpha_s$. The solid line stands for the case of $m_s = 0.5\text{GeV}$ and the dotted line for the heavier mass case of $m_s = 1.5\text{GeV}$. The difference between the upper and the lower frames lies in $\alpha_s$, where the former is for $\alpha_s = 0.02$ and the latter is for $\alpha_s = 0.08$.

evaluation of the integral in Eq.(1), we first found out the complex poles of the integrand, tracking their steps with parameters evolution. Then two opposite tendencies of the potential variation in terms of different parameters have been reported. The first tendency is to suppress the potential amplitude, slow down the oscillation "frequencies" and shift the peaks' centers to long distance. This tendency is caused by either decreasing $\alpha_s$ and $T$ or increasing $m$ and $m_s$. The second opposite tendency is to enhance the potential amplitude and frequencies, shifting the peaks center to short distance. As expected, this tendency is due to the increase of $\alpha_s$ and $T$ or decrease of $m$ and $m_s$. In brief, the potential between two fermions can be either long-weak or short-strong, according to the competition among the two groups of parameters. In some circumstances when the parameters are carefully "tuned", this potential is on the chance of being qualified strong and long enough to produce liquid state.

We would like to point out here that the behaviors of the potential are governed by the static modes which are determined by the dispersion relation of plasma. The Debye screening, which describes the high temperature environment, can be obtained in the HTL approximation to the boson self-energy where the external momenta are soft. This implies only soft modes at $q \sim gT$ get involved in the dielectric function when evaluating the interparticle potential. Therefore only those static modes in the soft region could be selected via dispersion relation and finally promise Debye screening. In other words, to reach Debye screening, one should take the high temperature limit by employing the HTL approximation from the beginning in the dielectric function to separate the soft modes, but not take the high temperature limit of $q_r$ and $q_i$ in Eq. (2) directly. As a matter of fact, the HTL static modes which are adapted to describe the extremely high temperature system may lose some significant details (hard mode contributions) in describing the medium at $1 \sim 3$ times critical temperature. Instead, the improved calculation in this paper picks up the the long distance oscillation tail of the potential which might be helpful to understand the low viscous mechanism of QGP with scalar bound state.

The interparticle potential is physically measurable at least in principle so we expect it is gauge independent. As to QED calculation, gauge problem does not emerge because the boson self-energy is gauge independent. While for QCD, the linear framework can only be persisted in the temporal axis gauge (TAG), but the one-loop calculations involving gluon self-coupling are various in different gauges even to the static modes. Anyway, we hope the discussion in TAG may give the qualitative features as well as correct directions of their evolution. More detailed investigations such as resummations are desirable in the future.
FIG. 6: Interparticle potential of QED and QED+sQED. The upper plot is at $T = 0.3\text{GeV}$ and the lower plot is at $T = 0.6\text{GeV}$. The dotted curves in both plots are for the pure QED potential and the solid curves are for the potential involving scalar bound states where the coupling strength $\alpha_s = 0.05$.

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[1] H. Sivak, A. Pérez and J. Diaz Alonso, Prog. Theor. Phys. 105, 961 (2001)
[2] J. Kapusta and T. Toimela, Phys. Rev. D 37, 3731 (1988)
[3] J. Diaz Alonso, A. Pérez and H. Sivak, Nucl. Phys. A 505, 695 (1989)
[4] J. Diaz Alonso, H. Sivak and A. Pérez, Phys. Rev. Lett. 73, 2536 (1994)
[5] A. L. Fetter and J. D. Walecka, Quantum Theory of Many Particle Systems (McGraw-Hill, New York, 1971)
[6] M. Le Bellac, Thermal Field Theory (Cambridge Univ. Press, Cambridge, 1996)
[7] R. Hosoya and K. Kajantie, Nucl. Phy. B 250, 666 (1985)
[8] S. Gavin, Nucl. Phys. A 435, 826 (1985)
[9] P. Danielewicz and M. Gyulassy, Phys. Rev. D 31, 53 (1985)
[10] D.W. von Oertzen, Phys. Lett. B 280, 103 (1992)
[11] G. Baym, H. Monien, C.J. Pethick and D.G. Ravenhall, Phy. Rev. Lett. B 64, 1867 (1990)
[12] H. Heiselberg, Phys. Rev. D 49, 4730 (1994).
[13] P. Arnold, G.D. Moore and L.G. Yaffe, JHEP 0011, 001 (2000)
[14] P. Arnold, G.D. Moore and L.G. Yaffe, JHEP 0305, 051 (2003)
[15] Liu Hui, Hou Defu and Li Jiarong, Eur. Phys. J. C 45, 459 (2006)
[16] S. Jeon, Phys. Rev. D 52, 3591 (1995).
[17] E. Wang, U. Heinz and X. Zhang, Phys. Rev. D 53,5978 (1996)
[18] E. Wang and U. Heinz, Phys. Lett. B 471, 208 (1999)
[19] M.H. Thoma, Phy. Lett. B 269, 144 (1991)
[20] H. Defu and L. Jiarong, Nucl. Phys. A 618, 371 (1997)
[21] G. Aarts and J. M. Martínez Resco, JHEP 0211, 022 (2002); ibid. 0402, 061 (2004); ibid. 05, 074 (2005)
[22] M.A. ValleBasagoiti, Phys. Rev. D 66, 045005 (2002)
[23] S. Jeon and L.G. Yaffe, Phys. Rev. D 53, 5799 (1996)
[24] M.E. Carrington, H. Defu and R. Kobes, Phys. Rev. D 62, 025010 (2000); ibid. 64, 025001 (2001)
[25] Liu Hui, Hou Defu and Li Jiarong, arXiv: hep-ph/0602221
[26] E. Shuryak, Nucl. Phys. A 750, 64 (2005)
[27] D. Teaney, Phys. Rev. C 68, 034913 (2003)
[28] A. Nakamura and S. Sakai, Phys. Rev. Lett. 94, 072305 (2005)
[29] M.H. Thoma, J. Phys. G 31, L7 (2005)
[30] A. Peshier and W. Cassing, Phys. Rev. Lett. 94, 172301 (1995)
[31] N.H. March and M.P. Tosi, Introduction to Liquid State Physics (World Scientific Publishing, Singapore 2002)
[32] P.A. Egelstaff, An Introduction to the Liquid State (Clarendon Press, Oxford 1992)
[33] E. Shuryak and I. Zahed, Phys. Rev. C 70, 021901(R) (2004)
[34] M. E. Peskin, and D. V. Schroeder, An Introduction to Quantum Field Theory (Addison-Wesley, MA 1995)
[35] J.I. Kapusta, Finite Temperature Field Theory (Cambridge Univ. Press, Cambridge, 1989)
[36] D.F. Litim and C. Manuel, Phys. Rev. D 64, 094013 (2001) Appendix
[37] S. Leupold and M.H. Thoma, Phys. Lett. B 465, 249 (1999)
[38] L.V. Keldysh, Zh. Eksp. Teor. Fiz. 47, 1515 (1964) [Sov. Phys. JETP 20, 1018 (1965)]
[39] M.E. Carrington, H. Defu, and M.H. Thoma, Eur. Phys. J. C7, 347 (1999)