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The Dissipative Potential Induced by QCD at Finite Temperature and Density[†]

by

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Abstract

In the framework of QCD at finite temperature we have obtained dissipative terms for the effective potential between q and $\bar{q}$ which would partly explain the J/ψ suppression in the Quark Gluon Plasma (QGP). The derivation of the dissipative potential for QGP is presented and the case for Hadron Matter (HM) is briefly discussed. The suppression effects are estimated based on simple approximations.

Key-words: J/ψ suppression: Quark Gluon Plasma.

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1. Introduction

It was supposed that a suppression of J/Ψ production in Relativistic Heavy Ion Collisions (RHIC) could be a clear signal for formation of a new matter phase, e.g. the Quark Gluon Plasma (QGP) [1], later the phenomenon was observed by the NA38 collaboration [2]. However, it is too early to celebrate the discovery of QGP, because rigorous studies indicate that in the Hadron Matter (HM) phase, the suppression phenomenon also exists. Thus many authors study the mechanisms which cause the suppression of J/Ψ production in the QGP and HM phases separately [1] [3]. They investigate how the mass of J/Ψ shifts with temperature and density of the QGP and HM medium, especially the dependence of the quark and gluon condensates on temperature and density are introduced to explain the mass shifts and then the suppression of J/Ψ production.

The RHIC experiments provide an atmosphere of hot and dense matter state, no matter it is in QGP or HM phases. The two situations are quite different in many aspects and an important research subject is to look for one or several clear signals to confirm or negate formation of QGP. The J/Ψ suppression is one of the candidates, but obviously, observation of such suppression is not enough for drawing a definite conclusion, one should investigate the details of the mechanisms which result in the suppression in QGP and HM. If one starts with a fundamental principle, i.e. QCD, instead of phenomenologically obtained potentials, he may expect to derive different potential forms for QGP and HM, and it just is the task of this work.

The effective potential between quark-antiquark or quark-quark can be written as

$$V(r) = -\frac{\alpha_{eff}}{r} + \sigma r. \tag{1}$$

In the regular zero-temperature theory $\alpha_{eff} = 4\alpha_s/3$ for meson and $2\alpha_s/3$ for baryon where $\alpha_s = g_s^2/4\pi$ and g_s is the QCD coupling constant. At finite temperature, the confinement constant becomes temperature-dependent [4],

$$\sigma(T) = \sigma(0)[\frac{T_{dec} - T}{T_{dec}}]^{\delta}\theta(T_{dec} - T), \tag{2}$$

where T_{dec} is the deconfinement temperature and δ is an uncertain parameter. Above T_{dec}, the linear confinement potential disappears, but it does not mean that the bound state dissolves, because a binding energy can also be provided by the Coulomb-type potential. By analysis of the medium state, a Debye screening mechanism is suggested as

$$V(r) = -\frac{\alpha_{eff}}{r} exp(-\frac{r}{r_D}), \tag{3}$$

where r_D is the Debye screening length. It is a function of temperature [1] and may be obtained by models. The dissolution condition can be written as

$$r_D \leq 0.84 r_B = 1.68/(\alpha_{eff} m_c).$$

Seeking for other possible mechanisms which may cause the J/Ψ suppression, we turn to re-study the situation. We will start with the QCD theory and its low energy phenomena at finite temperature and density.

Let us briefly retrospect how one derives the Coulomb-type potential from the field theory. Considering a t-channel scattering between two quarks or quark-antiquark, the amplitude in momentum space is

$$M = \overline{u}_1(\vec{p}_1)\gamma_\mu u_2(\vec{p}_2) G^{\mu\nu}(k^2) \overline{u}_3(\vec{p}_3)\gamma_\nu u_4(\vec{p}_4), \tag{4}$$

where $G^{\mu\nu}(k^2)$ is the full gluon propagator and for $q - \bar{q}$ interaction u_3 and u_4 become v_3 and v_4. There is also an extra s-channel annihilation diagram for $q - \bar{q}$, where u_2 and u_3 turn to v_2 and v_3. Setting $k_0 = 0$, a Fourier transformation of the amplitude gives rise to the effective potential in configuration space. Generally speaking, in $t-$channel the momentum transfer is space-like, i.e. $k^2 \leq 0$.

It is known that the photon and gluon propagators cannot provide an imaginary part, even at finite temperature and density in the common sense. The reason is that an absorptive part of the photon and gluon propagators corresponds to on-shell photon or gluon (gauge bosons), i.e. $k^2 = 0$, it indeed manifests an emission or absorption of a real photon or gluon from the quark (antiquark). At s-channel, the intermediate gauge bosons have momentum $k^2 = s = (p_q + p_{\bar{q}})^2 \geq 4m_q^2 > 0$, thus the $k^2 = 0$ condition cannot be satisfied, unless via loops. Whereas at t-channel the momentum transfer carried by the gauge boson is $k^2 = (p_q - p'_q)^2 \leq 0$, so it seems that an imaginary part can appear. But it is not true, because only at the boundary point $(p_q - p'_q)^2 = 0$, it means that if at both sides of the interaction, the quark is on its mass shell $E_q = E'_q$ and $\vec{p}_q = \vec{p}'_q$, i.e. the $\theta = 0$ forward scattering. For the Fourier transformation the integration should be restrained in a momentum conservation allowed region, so for the exact on-shell situation, the integration region is zero (see below in the context for details), namely $\int_0^0 |\vec{k}| f(k)$ would result in a null contribution to the absorptive part. By contraries, for the real part of the propagator which gives rise to the regular potential, $k^2 = 0$ condition is not imposed, so the integration region does not have any restriction (see below). That is why in the regular theories, photon and gluon propagators do not contribute an imaginary (absorptive) part, but only a real (dispersive) part.

However, in a bound state under consideration, the situation is different. If the heavy quark or antiquark (c, $\bar{c}$) is on-shell, it can be described as

$$p_\mu = m_c v_\mu, \tag{5}$$

where v_μ is the four-velocity, but in a bound state there is an integration region [5]

$$p_\mu = m_c v_\mu + k_\mu, \tag{6}$$

where k_μ is the "residue" momentum, and $|\vec{k}| \sim k_0 \leq \Lambda_{QCD}$. Thus the integration region allowed by the energy-momentum conservation is no longer zero, but can be from zero to Λ_{QCD} (or a smaller value than Λ_{QCD}). It can indeed contribute an absorptive part which finally results in a dissipative term in the effective potential. In the next section we will present the details.

Assuming the temperature and density are above T_{dec} and ρ_{dec}, the QGP phase is reached. Without the linear confinement term whose source is purely non-perturbative and obscure so far, we only need to take into account the one-gluon-exchange diagram which results in the interaction between quarks and quark-antiquark. It is well known that at $T = 0$, the leading term of the one-gluon-exchange contribution is the familiar Coulomb-type potential $-\alpha_{eff}/r$, but as the non-perturbative QCD effects which are characterized by the non-vanishing quark and gluon condensates are taken into account, the potential form is modified and corrections related to $1/m_Q$ (m_c or m_b) emerge [6][7]. When the temperature is non-zero and the density is above the regular one, more extra contributions appear. In this work, ignoring all spin and spin-orbit dependent terms and under some simple approximations, we obtain a new Coulomb-type potential as

$$V(r) = -\frac{\alpha_{eff}}{r}[1 + i(a + \alpha_s b)], \tag{7}$$

where a and b are functions of temperature and density. This new term turns the J/Ψ charmonium into a dissipative system and a dissolution is expected.

In the next section, we derive eq.(7) in the framework of QCD at finite temperature and in Sec.III, we discuss its significance and the situation for HM phase.

II. Formulation.

In the QGP atmosphere the propagator of fermion at finite temperature and density in momentum space can be written as [8]

$$S_F(k) = (\not{k} + m)[\frac{i}{k^2 - m^2} - 2\pi\delta(k^2 - m^2) f_F(k \cdot u)], \tag{8}$$

where

$$f_F(x) = \frac{\theta(x)}{e^{\beta(x-\mu)} + 1} + \frac{\theta(-x)}{e^{-\beta(x-\mu)} + 1}. \tag{9}$$

In the expression $\beta = 1/T$, u is the four-velocity of the medium, generally $u = (1,\ \vec{0})$ in the laboratory frame, so $k \cdot u = k_0$, μ is the chemical potential and is related to the density of the medium. Whereas the propagator of gluon reads

$$D_{\mu\nu}(k) = [\frac{i}{k^2} + 2\pi\delta(k^2) f_B(k \cdot u)](-g_{\mu\nu} + \frac{k_\mu k_\nu}{k^2}) \tag{10}$$

where

$$f_B(k\cdot u)=\frac{1}{e^{\beta k_0}-1}. \tag{11}$$

It is noted that we are working in the real-time scheme [9].

As aforementioned, if the charm-quark in J/Ψ bound state is strictly on mass shell, the t-channel scattering demands $k^2\le 0$, so that the fermion propagator does not have the temperature-dependent term because the intermediate fermion of non-zero mass cannot be an on-shell real particle. However, the soft gluon interaction can make the charm quark deviate from its mass shell, namely the k^2 can be equal to or greater than zero in eq.(6) and it represents an emission or absorption of real photon or gluon, ($k^2=0$ and $|\vec{k}|\sim\Lambda_{QCD}$), moreover a small time-like momentum transfer $k^2=m_q^2>0$ is also reasonable, then the intermediate fermion is an on-shell real particle.

Omitting the loop corrections, the lowest order Feynman diagrams are depicted in Fig.1. In general, the contribution can be decomposed into a form

$$f(\mu,T)=f_1(T)+g^2f_2(T)+g^2f_3(T,\mu), \tag{12}$$

where $f_1(T), f_2(T)$ and $f_3(T,\mu)$ are from one gluon exchange, gluon condensate and quark condensate diagrams respectively shown in Fig. 1. It is noted that only f_3 depends on the density of QGP, but not f_1 and f_2. Both $g^2f_2(T)$ and g The second term with the $\delta-$functions in the propagators correspond to the on-mass-shell intermediate particle, therefore turns to be a real particle, so it can feel the influence from the surrounding atmosphere, namely the density and temperature.

The traditional way for obtaining the potential is only to deal with the off-shell part of the propagator, namely to set $k_0=0$ and carry out a three-dimensional Fourier transformation of the two quark scattering amplitude. All the details including the loop corrections and non-perturbative effects are presented in literatures [6][7] [10][11], instead this work is only focused on the temperature and density-dependent parts which would contribute additional modification terms to the potential, we denote them as $V_G^T(r)$ and $V_q^T(r)$ corresponding to Fig.1 (a) and (b) respectively.

From Fig.1 (a), one has

$$V_G^T(r)=\frac{1}{M_{J/\Psi}}(\frac{-16\pi i}{3}\alpha_s)(-2\pi)\int\frac{d^4k}{(2\pi)^4}\delta(k^2)e^{-ik\cdot(x-y)}\frac{1}{e^{\beta k_0}-1}, \tag{13}$$

where the factor $(\frac{-16\pi i}{3}\alpha_s)$ coming as a common factor due to the color singlet condition of hadrons and it makes the leading Coulomb term be $-\frac{4\alpha_s}{3r}$ as required. It is also noted that here one does not need to set $k_0=0$ as in ref.[10] and the final integration region for $|\vec{k}|$ is from 0 to Λ_{QCD} as discussed in section I. The factor $\frac{1}{M_{J/\Psi}}$ guarantees the dimension of $V_G(r)$ right and comparing with the traditional way for deriving the potential, there is an extra integral over k_0, so a factor $\frac{1}{E}\approx\frac{1}{M_{J/\Psi}}$ is needed. This is equivalent to multiply a time factor τ ($\tau\cdot E\sim 1$).

Taking the spontaneous requirement $x_0=y_0$ (potential means an instantaneous interaction) a straightforward calculation gives

$$\begin{aligned}V_G^T(r) &= \frac{-16\pi i}{3}\alpha_s\frac{1}{(2\pi)^2}\cdot\frac{1}{rM_{J/\Psi}}\sum_{n=1}^{\infty}\frac{1}{n^2\beta^2+r^2}\times\\ &\quad [r(1-e^{-n\beta\Lambda}\cos\Lambda r)-n\beta e^{-n\beta\Lambda}\sin\Lambda r],\end{aligned} \tag{14}$$

where $\Lambda\le\Lambda_{QCD}$ is a parameter and in the expression we deliberately pull out the $1/r$ factor (see below). It is a series which absolutely converges for any finite r and β. If r is not very large (a few tenths of fm in our case), one only needs to take first a few terms for numerical computations. One can notice that if $\Lambda\sim 0$, i.e. for quarks are exactly on mass-shell, $V_G^T(r)$ vanishes and at $T\to 0$, it is also zero, this is consistent with our common knowledge. To see its meaning and avoiding tedious calculation, setting $\Lambda\to\infty$ we approximate eq.(14) to an integral as

$$\sum_{n=1}^{\infty}\frac{r}{n^2\beta^2+r^2}\sim\int_1^{\infty}dx\frac{r}{x^2\beta^2+r^2}=\frac{1}{\beta}(\frac{\pi}{2}-arctan\frac{\beta}{r}).$$

Thus the $V_G^T(r)$ is recast as

$$V_G^T(r) = \frac{-16\pi i}{3}\alpha_s \frac{1}{(2\pi)^2}\frac{1}{\beta r}(\frac{\pi}{2} - arctan\frac{\beta}{r}).$$

Below we will evaluate its contribution in terms of the series solution (13).

Similarly, for the quark contribution shown in Fig.1 (b), one has

$$\begin{aligned} V_q^T(r) &= \frac{-32\pi i}{9}\alpha_s^2 \frac{<\psi_q\overline{\psi}_q>_T}{m_q^3}\cdot\frac{1}{r}[\frac{1-\cos\mu r}{r} - \sum_{n=1}^{\infty}(-1)^n\frac{1}{n^2\beta^2+r^2}\times \\ & [2r\cos\mu r - (r\cos\Lambda r + n\beta\sin\Lambda r)e^{-n\beta\Lambda}(1+e^{\mu r})], \end{aligned} \quad (15)$$

for $\Lambda \geq \mu$ and

$$\begin{aligned} V_q^T(r) &= \frac{-32\pi i}{9}\alpha_s^2 \frac{<\psi_q\overline{\psi}_q>_T}{m_q^3}\cdot\frac{1}{r}[\frac{1-\cos\mu r}{r} - \sum_{n=1}^{\infty}(-1)^n\frac{e^{-n\beta\mu}}{n^2\beta^2+r^2}\times \\ & [r(2-(e^{n\beta\Lambda}-e^{-n\beta\Lambda})\cos\Lambda r) + n\beta\sin\Lambda r(e^{n\beta\Lambda}-e^{-n\beta\Lambda})]. \end{aligned} \quad (16)$$

for $\Lambda < \mu$, it is obvious that as $\Lambda \to 0$, $V_q^T(r)$ vanishes. While deriving eq.(15), we approximate $E = \sqrt{\vec{k}^2 + m_q^2} \approx |\vec{k}|$ in the integral due to the smallness of m_q. In the expression $<\psi_q\overline{\psi}_q>_T$ is the quark condensate at finite temperature T and its significance will be discussed in the next section.

III. Discussions.

The additional potential terms induced by QCD at finite temperature and density are imaginary and tend to zero as the temperature and density approach to zero. It makes sense because if there is no hot and dense atmosphere for the gluon and quarks, the additional terms do not exist at all, so this scenario would not affect the regular lifetime (about $7.5 \pm 0.4 \times 10^{-21}$ sec.) of the J/Ψ which mainly is determined by the s-channel annihilation process of $q\bar{q}$ into three gluons. When T is very high (a few hundred MeV) and $\rho \gg \rho_0$, the damping factor becomes substantial and it is the QGP situation. (see below)

The temperature and density effects turn the Hamiltonian of J/Ψ into complex, and the new Hamiltonian describes a dissipative quantum system. Because it is dissipative, the system would dissolve after a certain time, for example, into $D\overline{D}$ or $D_s\overline{D}_s$ etc. Dissolution to $D\overline{D}$ or $D_s\overline{D}_s$ needs to absorb energies from atmosphere because $M_{J/\Psi} < 2M_D$, in QGP case, it is possible, whereas in the vacuum circumstance, due to the kinematic constraint of the final state phase space, even there were a dissipative term, the dissolution would not occur.

From eqs.(14, 15) one can notice that the coefficients of i in $V_G^T(r)$ and $V_q^T(r)$ are always negative, This is a very important point because it makes the system dissipative.

Since $V_G^T(r)$ and $V_q^T(r)$ are very complicated functions of r, it would be extremely difficult to solve the Schrödinger equation with such a potential. Therefore to see the physical significance, we would choose a simple but reasonable approximation.

(a) Estimation of the dissipation effects.

(i) Looking at the additional Hamiltonian, one can note that they can be written as $if_{q(G)}(r)\frac{1}{r}$ where $f_{q(G)}(r)$ are complicated functions of r and given in eqs.(14, 15). Our approximation is to treat $f_{q(G)}(r)$ as average values instead of functions of r, namely, we approximate $f_{q(G)}(r)$ as $f_{q(G)}(\bar{r})$ where $\bar{r}$ is the average radius of J/Ψ. Matsui [1] suggests $0.2 \leq \bar{r}_{J/\Psi} \leq 0.5$ fm.

(ii) As $\sigma(T)$ disappears above the deconfinement temperature T_{dec}, the stationary Schrödinger equation becomes

$$-\frac{1}{2m_{red}}\nabla^2\phi - \frac{4\alpha_s}{3r}[1 + i(a + \alpha_s b)]\phi = E\phi \quad (17)$$

where ϕ is the wavefunction of J/Ψ, m_{red} is the reduced mass which equals $m_c/2$, and

$$a = \frac{1}{\pi}\sum_{n=1}^{\infty}\frac{1}{n^2\beta^2+\bar{r}^2}\times[\bar{r}(1 - e^{-n\beta\Lambda}\cos\Lambda\bar{r}) - n\beta e^{-n\beta\Lambda}\sin\Lambda\bar{r}]$$

$$\stackrel{\Lambda\to\infty}{\longrightarrow} \quad \frac{1}{\pi}\frac{1}{\beta}(\frac{\pi}{2} - arctan\frac{\beta}{\bar{r}}) \tag{18}$$

$$\begin{aligned} b &= \sum_{q=u,d,s}\frac{8}{3}\frac{<\psi_q\overline{\psi}_q>}{m_q^3}[\frac{1-\cos\mu\bar{r}}{\bar{r}} - \sum_{n=1}^{\infty}(-1)^n\frac{1}{n^2\beta^2+\bar{r}^2}\times \\ &\quad [2\bar{r}\cos\mu\bar{r} - (\bar{r}\cos\Lambda\bar{r} + n\beta\sin\Lambda\bar{r})e^{-n\beta\Lambda}(1+e^{\mu\bar{r}})], \qquad \text{for } \Lambda \geq \mu; \end{aligned} \tag{19}$$

$$\begin{aligned} b' &= \sum_{q=u,d,s}\frac{8}{3}\frac{<\psi_q\overline{\psi}_q>}{m_q^3}[\frac{1-\cos\mu\bar{r}}{\bar{r}} - \sum_{n=1}^{\infty}(-1)^n\frac{e^{-n\beta\mu}}{n^2\beta^2+\bar{r}^2}\times \\ &\quad [\bar{r}(2-(e^{n\beta\Lambda}+e^{-n\beta\Lambda})cos\Lambda\bar{r}) + n\beta\sin\Lambda\bar{r}(e^{n\beta\Lambda}-e^{-n\beta\Lambda})] \qquad \text{for } \Lambda < \mu. \end{aligned} \tag{20}$$

The eigenenergy of eq.(17) is simple and reads

$$(E_j)_{eff} = \frac{-8m_{red}}{9j^2}(\alpha_s)^2_{eff}, \tag{21}$$

where $(\alpha_s)_{eff} = \alpha_s[1+i(a-\alpha_s b)]$. For the ground state $j=1$.

(iii) The evolution of the quantum system can be expressed as

$$\begin{aligned} \phi_{J/\Psi}(t) &= \phi_{J/\Psi}(0)e^{-iE_{eff}t} \\ &= \phi_{J/\Psi}(0)exp(i\frac{8m_{red}\alpha_s^2}{9}[1-(a+\alpha_s b)^2]t)\cdot exp(\frac{-16m_{red}}{9}(a+\alpha_s b)\alpha_s^2 t). \end{aligned} \tag{22}$$

There exists a damping factor. In the calculations, we ignore the temperature dependence of α_s [12].

(b) Numerical evaluations.

(i) If the QCD expansion (including the condensates) converges, $|b\alpha_s|$ should be smaller than a. Typically, $\alpha_s \approx 0.3$ in the potential model, the term $\alpha_s b$ can compete with the term a. b is proportional to the quark condensate at finite temperature $<\psi_q\overline{\psi}_q>_T$ which decreases with the increase of temperature. However, we will argue in the following that the ratio $\frac{<\psi_q\overline{\psi}_q>_T}{m_q^3}$ where m_q is the constituent quark mass is almost independent of temperature.

The constituet quark mass is defined as the pole of the quark propagator

$$\Sigma(p^2 = m_q^2) = m_q, \tag{23}$$

where $\Sigma(p^2)$ is the dynamical mass related to quark condensate [16]

$$\Sigma(p^2) \sim \frac{1}{p^2} <\psi_q\overline{\psi}_q> \alpha_s(p^2). \tag{24}$$

Hence the ratio $\frac{<\psi_q\overline{\psi}_q>_T}{m_q^3}$ depends only on $\alpha_s(m_q^2)$ which dependence on temperature is ignored. Therefore, the parameter b is almost temperature independent.

(ii) Thus we have

$$|\phi_{J/\Psi}(t)|^2 = |\phi_{J/\Psi}(T=0)|^2 e^{(\frac{-2t}{\tau_0})}.$$

A typical time factor is about $\tau_0 \sim 7\times10^{-22}$ sec., for $0.2 \leq \bar{r} \leq 0.5$ fm and here we use $\bar{r} = 0.4$ fm. Even though this numerical value cannot be taken very seriously, the order of magnitude is reasonable.

The size of the collision region is about $10 \sim 100$ fm, a particle produced at the center of the region needs $1\times10^{-22} \sim 10^{-21}$ sec. to travel to the boundary, so τ_0 is of the same order as the traverse time and definitely results in a suppression of J/Ψ production (in fact a bulk decay) observed in experiments.

In fact, in such a small time interval, according to the regular theory c and $\bar{c}$ hardly combine into a bound state, so it is equivalent to the effective screening.

(iii) It is worth noting that only b depends on the density via the chemical potential μ. From eq.(15), numerical results show that as $|\mu|$ gets larger the contribution of $V_q^T(r)$ becomes more competitive to $V_G^T(r)$, but by the common sense, $|b\alpha_s| < a$.

(c) For the hadron matter phase, gluons and quarks do not directly feel the temperature and density of the hadron medium, therefore the gluon propagator cannot be influenced by the medium. In this case the temperature and density effects can appear in two ways. One is that the quark and gluon condensates are modified in the medium but different from that in QGP, [15], while another way is via loops in s-channel.

Our results show that the temperature and density effects can cause a suppression of J/Ψ production, but the mechanism is different from the traditional Debye screening effect, i.e. it contributes a dissipative part to the potential. This additional term makes the quantum system dissolve by a time scale 7×10^{-22} sec. This mechanism can only exist in QGP phase but not in the HM phase.

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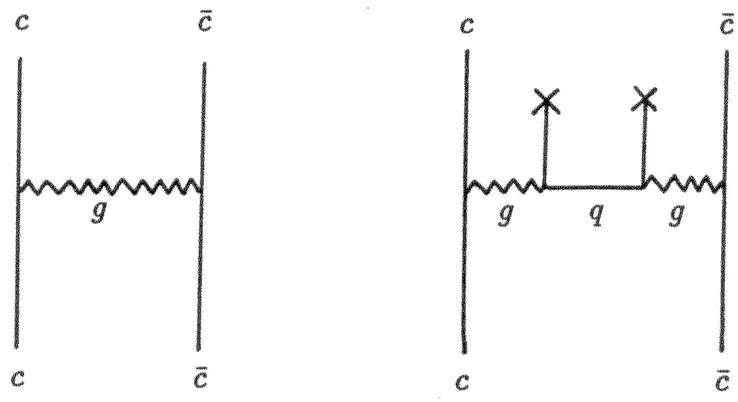
Figure Caption

Fig.1.The Feynman diagrams of t-channel $c - \bar{c}$ scattering where the charm-quark may deviate from its mass shell a bit.



(a) (b)

Fig.1

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