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Heavy hadron spectroscopy theory

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We discuss the multi-hadron systems containing a heavy hadron. We consider the heavy quark spin symmetry at leading order in the $1/m_{\rm Q}$ expansion with a heavy quark mass $m_{\rm Q}$. We discuss that there exist doublet/singlet states as degenerate/non-degenerate states in mass for any internal structure of the multi-hadron systems. As concrete examples, we adopt the heavy meson effective theory conserving the heavy quark spin, and apply it to the hadron molecule $\bar{P}^{(*)}N$ composed of the heavy meson $\bar{P}^{(*)} = (\bar{Q}q)_{\rm spin\,0(1)}$ (a heavy antiquark $\bar{Q}$ and a light quark q) and a nucleon N, which leads to the doublet state. We further show that $\bar{P}^{(*)}$ mesons are degenerate in mass in nuclear matter, and that the doublet state exists. Behind the mass degeneracy in the multi-hadron systems, a new structure called the spin-complex is proposed as a specific form of the brown muck. At next-to-leading order, we show that, in the heavy hadron mass formula, the terms conserving (breaking) the heavy quark spin symmetry are related to the interaction vertices of the chromoelectric (chromomagnetic) gluon and the heavy quark. By utilizing this property, we propose to use heavy hadrons as probes to investigate the modifications of the gluon fields in medium. Concretely, we perform the analysis of the in-medium masses of $\bar{P}^{(*)}$ meson in nuclear matter based on the heavy meson effective theory, and predict that the chromoelectric gluon is enhanced, while the chromomagnetic gluon is suppressed.

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1. Introduction

Recently, many exotic hadrons called X, Y, Z have been reported from experiments [1]. They have extraordinary masses, decay widths, decay patterns and so on, which are not expected from the conventional quark model. Though many theoretical interpretations, such as compact multi-quarks, hadronic molecules and so on, have been proposed, there are still many open questions. The problems are intimately related to the fundamental questions of the QCD.

One of the interesting points in the heavy quark systems is the existence of large mass scale given by the heavy quark mass m_{Q}. Then, the properties of the systems are described systematically by the $1/m_{\mathrm{Q}}$ expansion, which are given by the heavy quark effective theory (HQET). The leading order in the expansion is relevant to the heavy quark symmetry (HQS) [2, 3]. Thanks to the HQS, the heavy quark systems exhibit simple behaviors in mass spectrum in the heavy quark mass limit. They are the HQS doublet (singlet) having the mass (non-)degeneracy. The mass (non-)degeneracy must always exist irrespective to any quark and gluon configurations of the heavy systems. Therefore, the HQS doublet/singlet is a general concept which can be applied to, not only compact multi-quarks, but also to hadronic molecules as well as to heavy hadrons in nuclear matter and so on [4]. The next-to-leading order (NLO) terms in the $1/m_{\mathrm{Q}}$ expansion is related to the chromoelectric and chromomagnetic gluons coupled to a heavy quark. The NLO terms are included in the heavy hadron mass formula given by a series of $1/m_{\mathrm{Q}}$. By investigating the coefficients of the heavy hadron mass formula, we can obtain the information about the chromoelectric and chromomagnetic gluons in the heavy hadrons [5]. This property gives us a useful method to probe how the gluon fields are modified in medium at finite temperature and baryon number density.

The article is organized as followings. In Sec. 2, we give the formalism of the HQET. In Sec. 3, we consider the heavy quark limit, and introduce the new structure called the “spin-complex." In Sec. 4, we discuss the NLO in the $1/m_{\mathrm{Q}}$ expansion. Sec. 5 is devoted to the conclusion.

2. Heavy quark effective theory

In the HQET [2, 3], the four-momentum of the heavy quark Q is separated as $p^\mu = m_{\mathrm{Q}} v^\mu + k^\mu$ with the four-velocity v^μ ($v^2 = 1$) and the residual four-momentum k^μ whose scale is much smaller than m_{Q}. Then, we introduce the effective field for the heavy quark $Q_v(x) = e^{im_{\mathrm{Q}} v\cdot x} \frac{1+\not{v}}{2} Q(x)$ for the original heavy quark field $Q(x)$. The effective Lagrangian including $\mathscr{O}(1/m_{\mathrm{Q}})$ is given by

$$\mathscr{L}_{\mathrm{HQET}} = \overline{Q}_v v\cdot iD Q_v + \overline{Q}_v \frac{(iD_\perp)^2}{2m_{\mathrm{Q}}} Q_v - c(\mu) g_{\mathrm{s}} \overline{Q}_v \frac{\sigma_{\alpha\beta} G^{\alpha\beta}}{4m_{\mathrm{Q}}} Q_v + \mathscr{O}(1/m_{\mathrm{Q}}^2), \tag{2.1}$$

with $D_\perp^\mu = D^\mu - v^\mu\, v\cdot D$ for the covariant derivative $D^\mu = \partial^\mu + ig_{\mathrm{s}} A^\mu$ with the coupling constant g_{s} and the gluon field $A^\mu = A^{a\mu} T^a$ ($a = 1,\dots,8$). In the third term, the gluon field tensor $G^{\alpha\beta}$ is given by $[D^\alpha, D^\beta] = ig_{\mathrm{s}} G^{\alpha\beta}$, and the Wilson coefficient $c(\mu)$ is introduced from the matching to QCD at energy scale $\mu \simeq m_{\mathrm{Q}}$. The heavy quark symmetry (the heavy-flavor symmetry and the heavy-quark-spin symmetry) is conserved at $\mathscr{O}(1/m_{\mathrm{Q}}^0)$, while it is not generally conserved at $\mathscr{O}(1/m_{\mathrm{Q}})$. In the terms with $\mathscr{O}(1/m_{\mathrm{Q}})$ in Eq. (2.1), the first term breaks the heavy-flavor symmetry but still conserves the heavy-quark-spin symmetry, while the second term breaks both symmetries. The QCD Lagrangian for light quarks and gluons is unchanged.

Because the HQET Lagrangian (2.1) is given by a series of $1/m_{\mathrm{Q}}$, the mass of the heavy hadron is given also by by a series of $1/m_{\mathrm{Q}}$. This expansion is applied to the heavy hadrons, not only in vacuum, but also in medium at finite temperature T and/or baryon number density ρ, as long as the typical energy scale of T and ρ is much smaller than the heavy quark mass. The mass of the heavy hadron is given in the rest frame with $v_{\mathrm{r}}^{\mu} = (1, \vec{0})$ in the following form

$$M_{\mathrm{H}}(T,\rho) = m_{\mathrm{Q}} + \bar{\Lambda}(T,\rho) - \frac{\lambda_1(T,\rho)}{2m_{\mathrm{Q}}} + 4\vec{S}_{\mathrm{Q}}\cdot\vec{S}_{\mathrm{L}}\frac{\lambda_2(T,\rho;m_{\mathrm{Q}})}{2m_{\mathrm{Q}}} + \mathscr{O}(1/m_{\mathrm{Q}}^2). \tag{2.2}$$

with the matrix elements given as

$$\frac{1}{2M_{\mathrm{H}}}\langle \tilde{\mathrm{H}}_{v_{\mathrm{r}}}(T,\rho)|\frac{\beta(\alpha_{\mathrm{s}})}{4\alpha_{\mathrm{s}}}G^2|\tilde{\mathrm{H}}_{v_{\mathrm{r}}}(T,\rho)\rangle = \bar{\Lambda}(T,\rho), \tag{2.3}$$

$$\langle \mathrm{H}_{v_{\mathrm{r}}}(T,\rho)|\overline{Q}_{v_{\mathrm{r}}} g_{\mathrm{s}}\vec{x}\cdot\vec{E} Q_{v_{\mathrm{r}}}|\mathrm{H}_{v_{\mathrm{r}}}(T,\rho)\rangle = -\frac{\lambda_1(T,\rho)}{m_{\mathrm{Q}}}, \tag{2.4}$$

$$\frac{1}{2}c(\mu)\langle \mathrm{H}_{v_{\mathrm{r}}}(T,\rho)|\overline{Q}_{v_{\mathrm{r}}} g_{\mathrm{s}}\vec{\sigma}\cdot\vec{B} Q_{v_{\mathrm{r}}}|\mathrm{H}_{v_{\mathrm{r}}}(T,\rho)\rangle = 8\vec{S}_{\mathrm{Q}}\cdot\vec{S}_{\mathrm{L}}\lambda_2(T,\rho;m_{\mathrm{Q}}), \tag{2.5}$$

where $E^i = -G^{0i}$ is the chromoelectric gluon field and $B^i = \varepsilon^{ijk}G^{jk}$ is the chromomagnetic gluon field ($i,j,k = 1,2,3$), and the hadron state is denoted by $|\mathrm{H}_{v_{\mathrm{r}}}(T,\rho)\rangle$ in the medium at T and ρ. The first equation (2.3) originates from the scale anomaly in the trace of the energy-momentum tensor in QCD [6]. We introduce the Gell-Mann–Low function $\beta(\alpha_{\mathrm{s}}) = \mu \mathrm{d}\alpha_{\mathrm{s}}(\mu)/\mathrm{d}\mu$. The equation (2.4) is derived from the virial theorem as shown in Ref. [7]. Here $\vec{x}$ denotes the position of the center-of-mass of the system, which should coincide with the position of the heavy quark in the heavy quark limit. The equation (2.5) is straightforwardly obtained, because $\sigma_{\alpha\beta}$ contains the Pauli matrices $\vec{\sigma}$ for the heavy quark spin.

3. Heavy quark limit

3.1 Heavy quark spin and brown muck

Let us consider the heavy quark limit $m_{\mathrm{Q}} \to \infty$. This is an ideal case, but gives us essential features of the heavy quark systems. In this limit, only the first term in Eq. (2.1) survives. Importantly, it is invariant under the heavy quark spin transformation, namely $Q_v \to SQ_v$ with $S \in \mathrm{SU}(2)_{\mathrm{spin}}$. This is called the HQS[1]. Suppose that the total spin $\vec{J}$ of the hadron containing the heavy quark is decomposed into the sum of the heavy quark spin $\vec{S}$ and the total spin $\vec{j}$ of the remaining light components (light quarks and gluons) in the hadron; $\vec{J} = \vec{S} + \vec{j}$. Evidently $\vec{J}$ is a conserved quantity, and $\vec{S}$ is also conserved in the heavy quark limit. As a consequence, we find that $\vec{j}$ is also a conserved quantity. We should note that the carrier of $\vec{j}$ is a highly non-perturbative object composed of the light quark and gluons. Nevertheless, $\vec{j}$ is conserved as long as the heavy quark limit is taken. We call the object having $\vec{j}$ the brown muck, as an ensemble of the light quarks and gluons.

Thanks to the conservation of $\vec{j}$, there exist a pair of heavy hadrons which are degenerate in mass. They are two states with total spin $J_- = j - 1/2$ and $J_+ = j + 1/2$, respectively, with $j \geq 1/2$. We call them the HQS doublet. When $j = 0$, there exist only the state with $J_+ = 1/2$, and there

[1] In the present discussion, we concentrate only on the heavy quark spin, and we leave the heavy flavor symmetry.

is no partner corresponding to J_-. In this case, the state does not form a degenerate state. We call this the HQS singlet. The HQS doublet/singlet should exist in any heavy hadrons as long as the heavy quark mass limit is taken. In the real world, however, a heavy quark has a finite mass. Nevertheless, the mass degeneracy in the HQS doublet is seen approximately in charm and bottom hadrons. For example, the mass splitting between $\bar{D}$ and $\bar{D}^*$ mesons in charm is about 140 MeV, and that of B and B^* mesons in bottom is about 45 MeV. The small mass splittings are also seen in baryon sectors. In those examples, the possible configurations of the brown muck are given as a light quark q in heavy mesons and a diquark qq in heavy baryons from the constituent quark model.

3.2 Spin-complex

Let us consider hadronic molecules containing a heavy quark [4]. In hadronic molecules, the effective degrees of freedom can be given by hadrons, not by quarks and gluons. As a simple system, we may consider the hadronic molecule, the $\bar{P}^{(*)}N$ state composed of a $\bar{P} \sim (\bar{Q}q)_{\mathrm{spin}\,0}$ or $\bar{P}^* \sim (\bar{Q}q)_{\mathrm{spin}\,1}$ meson (a heavy antiquark $\bar{Q}$ and a light quark q) and a nucleon N. This is genuinely flavor exotic, because the minimal configuration $\bar{Q}qqqq$ cannot be reduced to that of the normal baryon[2]. Let us investigate how the HQS doublet/singlet is realized in the hadronic molecules.

To discuss concretely, we consider the $\bar{P}^{(*)}N$ states with $1/2^-$ and $3/2^-$ [8]. Their components are given by the basis sets with three channels, $\bar{P}N(^2S_{1/2})$, $\bar{P}^*N(^2S_{1/2})$ and $\bar{P}^*N(^4D_{1/2})$, for $1/2^-$ and by those with four channels, $\bar{P}N(^2D_{3/2})$, $\bar{P}^*N(^4S_{3/2})$, $\bar{P}^*N(^4D_{3/2})$ and $\bar{P}^*N(^2D_{3/2})$, for $3/2^-$. As an interaction between $\bar{P}^{(*)}$ and N, we adopt the one-pion exchange potential (OPEP) as discussed in Refs. [8]. We obtain the potentials in matrix forms with the basis sets; the 3×3 matrix $H_{1/2^-}$ for $1/2^-$ and the 4×4 matrix $H_{3/2^-}$ for $3/2^-$ (see Refs. [8] for the details). Note the tensor potentials in the off-diagonal components in the matrices. They indeed give a mixing between $\bar{P}N$ and $\bar{P}^*N$ channels with different angular momenta, L and $L\pm 2$, and hence provide a strong attraction leading to the formation of the bound/resonant states.

Let us discuss the HQS doublet/singlet in the present systems. Are there indeed any HQS doublet/or singlet state? How is the HQS doublet/singlet formed? We adopt unitary transformations for $H_{1/2^-}$ and $H_{3/2^-}$, and obtain the partly diagonalized matrices;

$$U_{1/2^-}^{-1} H_{1/2^-} U_{1/2^-} = \mathrm{diag}(H_{1/2^-}^{(0^+)}, H_{1/2^-}^{(1^+)}), \quad U_{3/2^-}^{-1} H_{3/2^-} U_{3/2^-} = \mathrm{diag}(H_{3/2^-}^{(1^+)}, H_{3/2^-}^{(2^+)}), \tag{3.1}$$

with unitary matrices $U_{1/2^-}$ and $U_{3/2^-}$, respectively. Interestingly, we find that the eigenvalues of $H_{1/2^-}^{(1^+)}$ are identical to those of $H_{3/2^-}^{(1^+)}$. That is, the mass of $1/2^-$ is the same as that of $3/2^-$. The mass degeneracy is not accidental, but is a consequence from the HQS. In the OPEP in our example, the same coupling constants for $\bar{P}\bar{P}^*\pi$ and for $\bar{P}^*\bar{P}^*\pi$ vertices leads the HQS in the systems.

In the previous section, we have discussed that there exist the HQS doublets/singlets for any configuration of the brown muck. In the present example of $\bar{P}^{(*)}N$, the brown muck is composed of (i) a light quark q in $\bar{P}^{(*)}$ meson, (iii) a nucleon N and (iii) the energy from the relative angular momentum between $\bar{P}^{(*)}$ and N. It is important to note that the configuration is different from the conventional picture for the brown mucks, e.g. a light quark q in heavy mesons and a pair of quarks (diquark) qq in heavy baryons in the constituent quark model. It is rather a complex system

[2]For $P^{(*)}N$ states with a $P^{(*)}$ meson (a anti-particle of $P^{(*)}$), we need to discuss the mixing with Qqq states.

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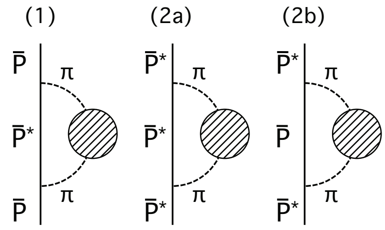

Figure 1: The diagram for the self-energy of $\bar{P}^{(*)}$ meson nuclear matter. The solid (dashed) line is a propagator for $\bar{P}^{(*)}$ (π) meson. The blob stands for the dressing by nucleon-hole pairs.

composed of *a light quark q* and *a nucleon N*. We call it the light spin-complex (or spin-complex in short), because it has the total spin as the conserved quantity and has a complex structure with light degrees of freedom of quarks and hadrons. In the present case, the total spin and parity of the spin-complex is 1^+, because the $1/2^-$ and $3/2^-$ states in $\bar{P}^{(*)}N$ are degenerate. Note that the spin-complex is defined when the Fock space is given (q and N in the present case). Nevertheless, the spin-complex is a useful object for analysis of the structure of the heavy quark systems.

3.3 Spin-complex in nuclear matter

Let us consider another example of the spin-complex for a $\bar{P}^{(*)}$ meson [4]. Whenever there is an attractive force between $\bar{P}^{(*)}$ and N, the $\bar{P}^{(*)}$ meson can be bound in nuclear matter with infinite volume [5, 10]. According to the general discussion in the previous subsection, there should exist the HQS doublet/singlet for a $\bar{P}^{(*)}$ meson in nuclear matter in the heavy quark limit. We will show concretely the presence of the HQS doublet.

We again consider the pion exchange between $\bar{P}^{(*)}$ and nucleons in nuclear matter, and analyze the lowest order in the perturbation for the interaction as shown in Fig. 1. The in-medium self-energies are given as $-i\Sigma_{\bar{P}}$ for the $\bar{P}$ meson in panel (1), and $-i\Sigma_{\bar{P}^*} = -i\Sigma_{\bar{P}^*}^{(\bar{P}^*)} - i\Sigma_{\bar{P}^*}^{(\bar{P})}$ with $\Sigma_{\bar{P}^*}^{(\bar{P}^*)} = (2/3)\Sigma_{\bar{P}}$ and $\Sigma_{\bar{P}^*}^{(\bar{P})} = (1/3)\Sigma_{\bar{P}}$ for the $\bar{P}^*$ meson in panels (2a,b). Two facts are important. First, thanks to the HQS, the same coupling strengths for the $\bar{P}\bar{P}^*\pi$ and $\bar{P}^*\bar{P}^*\pi$ vertices should be adopted. Second, both $\bar{P}$ and $\bar{P}^*$ with same mass should be considered simultaneously as the intermediate (virtual) states, as long as the states are allowed to exist. Then, we find that the contributions from (1) and (2a,b) in Fig. 1 are the same except for the overall coefficients. Their fraction 3 : 2 : 1 is understood intuitively from the counting of the spin degrees of freedom in the intermediate states. Consequently, we find that the self-energy of the P meson is equal to that of the $\bar{P}^*$ meson; $-i\Sigma_{\bar{P}} = -i\Sigma_{\bar{P}^*}$. Thus, we confirm the HQS doublet for a $\bar{P}^{(*)}$ meson in nuclear matter. The spin-complex is composed of a light quark q in $\bar{P}^{(*)}$ and multiple number of nucleon-hole pairs $NN^{-1} + NNN^{-1}N^{-1} + \ldots$ denoted by blobs in Fig. 1. The total spin and parity of the spin-complex is 1^+.

Recently the three-body system $\bar{P}^{(*)}NN$ with baryon number two has been investigated, and has been shown to be the HQS doublet in the heavy quark mass limit [9]. There, the spin-complex (qNN from q from $\bar{P}^{(*)}$ and two nucleons NN) has the total spin and parity $1/2^+$. It will be generally extended to systems with any baryon number, leading to the HQS doublet in the nuclear matter.

4. Probing gluon fields by heavy hadron — NLO in $1/m_{\mathrm{Q}}$ expansion —

In the previous section, we have considered the heavy quark limit and found that there exist the HQS doublet/singlet states accompanying the brown muck. Now let us consider the higher order of the HQET, namely the NLO in the $1/m_{\mathrm{Q}}$ expansion. As shown in Eqs. (2.4) and (2.5), $\lambda_1(T,\rho)$ and $\lambda_2(T,\rho;m_{\mathrm{Q}})$ $(Q=c,b)$ are related to the chromoelectric and chromomagnetic gluon fields interacting with the heavy quark. It means that the analysis of the $1/m_{\mathrm{Q}}$ expansion in the mass of the heavy hadrons are useful to probe the dynamics of the gluon fields. We will omit T, because it is an irrelevant variable in the present discussion.

In the nuclear matter at zero temperature and finite baryon number density ρ, we consider the $\bar{P}^{(*)}$ meson as a probe for medium modifications of gluon fields. To obtain numerical predictions for $\lambda_1(\rho)$ and $\lambda_2(\rho;m_{\mathrm{Q}})$, we consider the diagrams in Fig. 1, and adopt the approximation of the one-nucleon loop in the blobs. In those diagrams, the effects from NLO terms including the breaking of the HQS appear in the mass splitting between $\bar{P}$ and $\bar{P}^*$ mesons, as well as in the difference of the couplings for $\bar{P}\bar{P}^*\pi$ and $\bar{P}^*\bar{P}^*\pi$ vertices. To include the NLO terms systematically, we follow the invariance under the velocity-rearrangement ($v^\mu \to w^\mu + q^\mu/M$ with q^μ being a small momentum and M being the averaged mass of $\bar{P}$ and $\bar{P}^*$ mesons) [11], and obtain the Lagrangian for the heavy meson effective theory which has been rederived also in Ref. [5]. The breaking of the HQS is seen in the difference in $\bar{P}$ and $\bar{P}^*$ masses, and in the difference in the $\bar{P}\bar{P}^*\pi$ and $\bar{P}^*\bar{P}^*\pi$ vertices.

By setting the momentum cutoff parameter in the nucleon loop, we obtain the $\bar{P}^{(*)}$ self-energy in nuclear matter [5] (see also Ref. [10]), and from the mass formula (2.2) we obtain $\lambda_1(\rho)$ and $\lambda_2(\rho;m_{\mathrm{Q}})$ as functions of baryon number density ρ (Fig. 2.). We plot $\lambda_1(\rho)$ and $\lambda_2(\rho;m_{\mathrm{Q}})$ by dividing them by the values $\lambda_1(\rho=0)$ and $\lambda_2(\rho=0;m_{\mathrm{Q}})$ in vacuum. We find that $\lambda_1(\rho)$ increases as ρ increases, while $\lambda_2(\rho;m_{\mathrm{Q}})$ decreases. As $\lambda_1(\rho)$ and $\lambda_2(\rho;m_{\mathrm{Q}})$ are related to the chromoelectric and chromomagnetic gluons in medium (Eqs. (2.4) and (2.5)), respectively, we may interpret our result so that the chromoelectric (chromomagnetic) gluon field is enhanced (suppressed) in the nuclear matter in comparison with those in vacuum. The suppression of chromomagnetic gluon fields in nuclear matter indicates that the mass splitting between $\bar{P}$ and $\bar{P}^*$ mesons in nuclear medium are reduced than that in vacuum. This is consistent with the result from the QCD sum rule [12].

5. Summary and conclusion

We have discussed the multi-hadron systems containing a heavy quark. Based on the HQET, we have shown that the HQS doublet/singlet states exist in multi-hadron systems in the heavy quark limit. We have also shown that the $\bar{P}^{(*)}$ meson in nuclear matter belongs to the HQS doublet. At NLO in the $1/m_{\mathrm{Q}}$ expansion, we have discussed that the terms conserving (breaking) the HQS is sensible to the chromoelectric (chromomagnetic) gluon fields in the heavy hadron. As a concrete example, we have analyzed the in-medium masses of $\bar{P}^{(*)}$ mesons in nuclear matter, and shown that the chromoelectric gluons are enhanced, while the chromomagnetic gluons are suppressed.

As we have emphasized, the existence of the HQS doublet/singlet states in the heavy quark systems are independent of the specific structures of the systems. They will be useful to study many exotic heavy hadrons. To use heavy hadrons as probes of gluon fields is also applied to any medium at finite temperature, and even in the deconfinement phase such as the quark gluon plasma and the

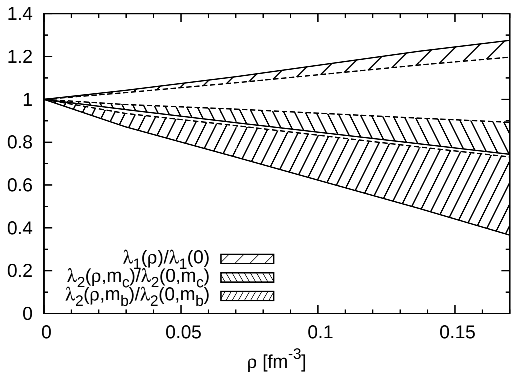

Figure 2: $\lambda_1(\rho)$ and $\lambda_2(\rho;m_Q)$ $(Q=c,b)$ for $\bar{P}^{(*)}$ meson as functions of baryon number density ρ in nuclear matter. The solid and dashed lines indicate the results for $(g, g_1/M_D, g_2/M_D) = (0.5, 0, -0.07)$ and $(0.4, 0, -0.17)$, respectively, with D meson mass M_D.

color superconductivity. Such subjects will be studied in experiments in KEK-Belle, J-PARC, GSI-FAIR, BNL-RHIC, CERN-LHC and so on.

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