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Research Article

Hidden-Beauty Broad Resonance $Y_b(10890)$ in Thermal QCD

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In this work, the mass and pole residue of resonance Y_b is studied by using QCD sum rules approach at finite temperature. Resonance Y_b is described by a diquark-antidiquark tetraquark current, and contributions to operator product expansion are calculated by including QCD condensates up to dimension six. Temperature dependencies of the mass m_{Y_b} and the pole residue λ_{Y_b} are investigated. It is seen that near a critical temperature ($T_c \simeq 190$ MeV), the values of m_{Y_b} and λ_{Y_b} decrease to 87% and to 44% of their values at vacuum.

1. Introduction

Heavy quarkonia systems provide a unique laboratory to search the interplay between perturbative and nonperturbative effects of QCD. They are nonrelativistic systems in which low energy QCD can be investigated via their energy levels, widths, and transition amplitudes [1]. Among these heavy quarkonia states, vector charmonium and bottomonium sectors are experimentally studied very well, since they can be detected directly in e^+e^- annihilations. In the past decade, observation of a large number of bottomonium-like states in several experiments increased the interest in these structures [2–6]. However, these observed states could not be conveniently explained by the simple $q\overline{q}$ picture of mesons. The presumption of hadrons containing quarks more than the standard quark content ($q\overline{q}$ or qqq) is introduced by a perceptible model for diquarks plus antidiquarks, which was developed by Jaffe in 1976 [7]. Later Maiani, Polosa, and their collaborators proposed that the X, Y, Z mesons are tetraquark systems, in which the diquark-antidiquark pairs are bound together by the QCD color forces [8]. In this color configuration, diquarks can play a fundamental role in hadron spectroscopy. Thus, probing the multiquark matter has been an intensely intriguing research topic in the past twenty years and it may provide significant clues to understand the nonperturbative behavior of QCD.

In 2007, Belle reported the first evidence of $e^+e^- \longrightarrow \Upsilon(1S)\pi^+\pi^-, \Upsilon(2S)\pi^+\pi^-$ and first observation for $e^+e^- \longrightarrow \Upsilon(3S)\pi^+\pi^-, \Upsilon(1S)K^+K^-$ decays near the peak of the $\Upsilon(5S)$ state at $\sqrt{s} = 10.87$ GeV [2]. Assigning these signals to $\Upsilon(5S)$, the partial widths of decays $\Upsilon(5S) \longrightarrow \Upsilon(1S)\pi^+\pi^-$ and $\Upsilon(5S) \longrightarrow \Upsilon(2S)\pi^+\pi^-$ were measured unusually larger (more than two orders of magnitude) than formerly measured decay widths of $\Upsilon(nS)$ states. Following these unusually large partial width measurement, Belle measured the cross sections of $e^+e^- \longrightarrow \Upsilon(1S)\pi^+\pi^-, \Upsilon(2S)\pi^+\pi^-$ and $\Upsilon(1S)\pi^+\pi^-$ and reported that the resonance observed via these decays does not agree with conventional $\Upsilon(5S)$ line shape. These observations led to the proposal of existence of new exotic hidden-beauty state analogous to broad $Y(4260)$ resonance in the charmonium sector, which is a Breit-Wigner shaped resonance with mass $(10888.4^{+2.7}_{-2.6} \pm 1.2)$ MeV/c^2, and width $(30.7^{+8.3}_{-7.0} \pm 3.1)$ MeV/c^2, and is called $Y_b(10890)$ [5]. In literature, there are several approaches to investigate the structure of exotic Y_b resonance. In [9], Y_b is considered as a $\Lambda_b\overline{\Lambda}_b$ bound state with a highly large binding energy. In [10, 11], Y_b is interpreted as a tetraquark and its mass is estimated by using QCD sum rules at vacuum.

Moreover, it is likely that at very high temperatures within the first microseconds following the Big Bang, quarks and gluons existed freely in a homogenous medium called the quark-gluon plasma (QGP). One of the first quark-gluon plasma signals proposed in the literature is suppression of J/ψ particles [12]. In 2011, CMS collaboration reported that charmonium states $\phi(2S)$ and J/ψ melt or were suppressed due to interacting with the hot nuclear matter created in heavy-ion interactions [13, 14]. Following these observations in the charmonium sector, CMS collaboration also reported suppression of bottomonium states, $\Upsilon(2S)$ and $\Upsilon(3S)$ relative to the $\Upsilon(1S)$ ground state [15, 16]. The dissociation temperatures for the Υ states are expected to be related to their binding energies and are predicted to be $2T_c, 1.2T_c$ and T_c for the $\Upsilon(1S), \Upsilon(2S)$, and $\Upsilon(3S)$ mesons, respectively, where T_c is the critical temperature for deconfinement [16–18].

In addition to these studies, one of the first investigations on thermal properties of exotic mesons was done in [19], in which the authors investigated in medium properties of $X(3872)$ under the hypothesis that it is a 1^{++} state or 2^{-+} state, but making no assumption on its structure. They estimated that the mass of the 1^{++} molecular state decreases with increasing temperature; the mass of charmonium or tetraquark state is almost stable.

Inspired by these findings and motivated by the aforementioned discussions, we focus on the Y_b resonance and its thermal behavior. This paper is organized as follows. In Section 2, theoretical framework of Thermal QCD sum rules (TQCDSR) and its application to Y_b are presented, and obtained analytical expressions of the mass and pole residue of Y_b are given up to dimension six operators. Numerical analysis is performed and results are obtained in Section 3. Concluding remarks are discussed in Section 4. The explicit forms of the spectral densities are written in Appendix.

2. Finite Temperature Sum Rules for Tetraquark Assignment

QCD sum rules (QCDSR) approach is based on Wilson's operator product expansion (OPE) which was adapted by Shifman, Vainshtein, and Zakharov [24] and applied with remarkable success to estimate a large variety of properties of all low-lying hadronic states [25–27]. Later, this model is extended to its thermal version that is firstly proposed by Bochkarev and Shaposnikov and led to many successful applications in QCD at $T \neq 0$ [28–35]. Very recently, in [36], the authors claimed that any QCDSR study on tetraquark states should contain $O(\alpha_s^2)$ contributions to OPE, which are unknown. It is a very important and strong argument for studying tetraquark states within QCDSR; however, those terms and their thermal behaviors are not known, and their calculation is beyond the scope of this work. Thus we follow the traditional sum rules to investigate the thermal behavior of hadronic parameters of Y_b, which were used very successfully in predicting properties of tetraquark states as well. In this work, we proceed with traditional sum rules analysis by using a tetraquark current, following several successful applications to exotic hadrons [37–42].

In this section, the mass and pole residue of the exotic Y_b resonance are studied by interpreting it as a bound $[bs][\overline{bs}]$ tetraquark via TQCDSR technique which starts with the two point correlation function

$$\Pi_{\mu\nu}(q,T) = i\int d^4x e^{iq\cdot x}\left\langle \Psi \left| \mathcal{T}\left\{\eta_\mu(x)\eta_\nu^\dagger(0)\right\}\right| \Psi\right\rangle, \quad (1)$$

where Ψ represents the hot medium state, $\eta_\mu(x)$ is the interpolating current of the Y_b state, and $\mathcal{T}$ denotes the time ordered product. The thermal average of any operator $\widehat{O}$ in thermal equilibrium is given as

$$\left\langle \widehat{O}\right\rangle = \frac{Tr\left(e^{-\beta\mathcal{H}}\widehat{O}\right)}{Tr\left(e^{-\beta\mathcal{H}}\right)}, \quad (2)$$

where $\mathcal{H}$ is the QCD Hamiltonian, and $\beta = 1/T$ is inverse of the temperature, and T is the temperature of the heat bath. Chosen current $\eta_\mu(x)$ must contain all the information of the related meson, like quantum numbers, quark contents and so on. In the diquark-antidiquark picture, tetraquark current interpreting Y_b can be chosen as [10]

$$\begin{aligned}\eta_\mu(x) = \frac{i\epsilon\tilde{\epsilon}}{\sqrt{2}}&\left\{\left[s_a^T(x)C\gamma_5Q_b(x)\right]\left[\bar{s}_d(x)\gamma_\mu\gamma_5C\overline{Q}_e^T(x)\right]\right.\\ &\left.+\left[s_a^T(x)C\gamma_5\gamma_\mu Q_b(x)\right]\left[\bar{s}_d(x)\gamma_5C\overline{Q}_e^T(x)\right]\right\},\end{aligned} \quad (3)$$

where $Q = b$, C is the charge conjugation matrix and a, b, c, d, e are color indices. Shorthand notations $\epsilon = \epsilon_{abc}$ and $\tilde{\epsilon} = \epsilon_{dec}$ are also employed in (3).

In TQCDSR, the correlation function given in (1) is calculated twice, as in two different regions corresponding two perspectives, namely the physical side (or phenomenological side) and the QCD side (or OPE side). By equating these two approaches, the sum rules for the hadronic properties of the exotic state under investigation are achieved. To derive mass and pole residue via TQCDSR, the correlation function is calculated in terms of hadronic degrees of freedoms in the physical side. A complete set of intermediate physical states possessing the same quantum number as the interpolating current are inserted into (1), and integral over x is handled. After these manipulations, the correlation function is obtained as

$$\begin{aligned}\Pi_{\mu\nu}^{\text{Phys}}(q,T) = &\frac{\left\langle \Psi\left|\eta_\mu\right| Y_b(q)\right\rangle_T\left\langle Y_b(q)\left|\eta_\nu^\dagger\right|\Psi\right\rangle_T}{m_{Y_b}^2(T)-q^2}\\ &+ \textit{subtracted terms},\end{aligned} \quad (4)$$

here $m_{Y_b}(T)$ is the temperature-dependent mass of Y_b meson. Temperature-dependent pole residue $\lambda_{Y_b}(T)$ is defined in terms of matrix element as

$$\left\langle \Psi\left|\eta_\mu\right| Y_b(q)\right\rangle_T = \lambda_{Y_b}(T)\, m_{Y_b}(T)\,\varepsilon_\mu, \quad (5)$$

where ε_μ is the polarization vector of the Y_b satisfying

$$\varepsilon_\mu\varepsilon_\nu^* = -g_{\mu\nu} + \frac{q_\mu q_\nu}{m_{Y_b}^2(T)}. \quad (6)$$

After employing polarization relations, the correlation function is written in terms of Lorentz structures in the form

$$\Pi_{\mu\nu}^{\text{Phys}}(q,T) = \frac{m_{Y_b}^2(T)\,\lambda_{Y_b}^2(T)}{m_{Y_b}^2(T) - q^2}\left(-g_{\mu\nu} + \frac{q_\mu q_\nu}{m_{Y_b}^2(T)}\right) + \dots, \tag{7}$$

where dots denote the contributions coming from the continuum and higher states. To obtain the sum rules, coefficient of any Lorentz structure can be used. In this work, coefficients of $g_{\mu\nu}$ are chosen to construct the sum rules and the standard Borel transformation with respect to q^2 is applied to suppress the unwanted contributions. The final form of the physical side is obtained as

$$\mathscr{B}\left(q^2\right)\Pi^{\text{Phys}}(q,T) = m_{Y_b}^2(T)\,\lambda_{Y_b}^2(T)\,e^{-m_{Y_b}^2(T)/M^2}, \tag{8}$$

here M^2 is the Borel mass parameter. In the QCD side, $\Pi_{\mu\nu}^{\text{QCD}}(q,T)$ is calculated in terms of quark-gluon degrees of freedom and can be separated into two parts over the Lorentz structures as

$$\Pi_{\mu\nu}^{\text{QCD}}(q,T) = \Pi_S^{\text{QCD}}\left(q^2,T\right)\frac{q_\mu q_\nu}{q^2} + \Pi_V^{\text{QCD}}\left(q^2,T\right)\left(-g_{\mu\nu} + \frac{q_\mu q_\nu}{q^2}\right), \tag{9}$$

where $\Pi_S^{\text{QCD}}(q^2,T)$ and $\Pi_V^{\text{QCD}}(q^2,T)$ are invariant functions connected with the scalar and vector currents, respectively. In the rest framework of Y_b ($\mathbf{q} = 0$), $\Pi_V^{\text{QCD}}(q_0^2,T)$ can be expressed as a dispersion integral,

$$\Pi_V^{\text{QCD}}\left(q_0^2,T\right) = \int_{4(m_b+m_s)^2}^{s_0(T)} \frac{\rho^{\text{QCD}}(s,T)}{s - q_0^2}\,ds + \cdots, \tag{10}$$

where corresponding spectral density is described as

$$\rho^{\text{QCD}}(s,T) = \frac{1}{\pi}\,\text{Im}\,\Pi_V^{\text{QCD}}(s,T). \tag{11}$$

The spectral density can be separated in terms of operator dimensions as

$$\rho^{\text{QCD}}(s,T) = \rho^{\text{pert.}}(s) + \rho^{\langle\bar{q}q\rangle}(s,T) + \rho^{\langle G^2\rangle + \langle\Theta_{00}\rangle}(s,T) + \rho^{\langle\bar{q}Gq\rangle}(s,T) + \rho^{\langle\bar{q}q\rangle^2}(s,T). \tag{12}$$

In order to obtain the expressions of these spectral density terms, the current expression given in (3) is inserted into the correlation function given in (1) and then the heavy and light quark fields are contracted, and the correlation function is written in terms of quark propagators as

$$\begin{aligned}
\Pi_{\mu\nu}^{\text{QCD}}(q,T) = {} & -\frac{i}{2}\int d^4x\,e^{iq\cdot x} \\
& \cdot \epsilon\tilde{\epsilon}\epsilon'\tilde{\epsilon}'\Big\langle\Big\{\text{Tr}\left[\gamma_\mu\gamma_5\widetilde{S}_b^{aa'}(-x)\,\gamma_5\gamma_\nu S_s^{bb'}(-x)\right] \\
& \times \text{Tr}\left[\gamma_5\widetilde{S}_s^{dd'}(x)\,\gamma_5 S_b^{ee'}(x)\right] \\
& + \text{Tr}\left[\gamma_5\widetilde{S}_b^{aa'}(-x)\,\gamma_5 S_s^{bb'}(-x)\,\gamma_\mu\right] \\
& \times \text{Tr}\left[\gamma_5\widetilde{S}_s^{dd'}(x)\,\gamma_5\widetilde{S}_b^{ee'}(x)\,\gamma_\nu\gamma_5 S_b^{bb'}(x)\right] \\
& + \text{Tr}\left[\gamma_5\widetilde{S}_b^{aa'}(-x)\,\gamma_5\gamma_\nu \times S_s^{bb'}(-x)\right] \\
& \cdot \text{Tr}\left[\gamma_5\widetilde{S}_s^{dd'}(x)\,\gamma_5\gamma_\mu S_b^{ee'}(x)\right] \\
& + \text{Tr}\left[\gamma_5\widetilde{S}_b^{aa'}(-x)\,\gamma_5 \times S_s^{bb'}(-x)\right] \\
& \cdot \text{Tr}\left[\gamma_5\widetilde{S}_s^{dd'}(-x)\,\gamma_5 S_b^{ee'}(x)\,\gamma_\nu\right]\Big\}\Big\rangle_T,
\end{aligned} \tag{13}$$

where $S_{s,b}^{ij}(x)$ are the full quark propagators and $\widetilde{S}_{s,b}^{ij}(x) = CS_{s,b}^{ijT}(x)C$ is used. The quark propagators are given in terms of the quark and gluon condensates [25]. At finite temperatures, additional operators arise due to the breaking of Lorentz invariance by the choice of thermal rest frame. Thus, the residual $O(3)$ invariance brings additional operators to the quark propagator at finite temperature. The expected behavior of the thermal averages of these new operators is opposite of those of the Lorentz invariant old ones [43]. The heavy-quark propagator in coordinate space can be expressed as

$$\begin{aligned}
S_b^{ij}(x) = {} & i\int\frac{d^4k}{(2\pi)^4}e^{-ik\cdot x}\left[\frac{\delta_{ij}(\not{k} + m_b)}{k^2 - m_b^2} - \frac{gG_{ij}^{\alpha\beta}}{4}\frac{\sigma_{\alpha\beta}(\not{k} + m_b) + (\not{k} + m_b)\sigma_{\alpha\beta}}{(k^2 - m_b^2)^2}\right. \\
& \left. + \frac{g^2}{12}G_{\alpha\beta}^A G_A^{\alpha\beta}\delta_{ij}m_b\frac{k^2 + m_b\not{k}}{(k^2 - m_b^2)^4} + \dots\right],
\end{aligned} \tag{14}$$

and the thermal light quark propagator is chosen as

$$\begin{aligned}
S_s^{ij}(x) = {} & i\frac{\not{x}}{2\pi^2x^4}\delta_{ij} - \frac{m_s}{4\pi^2x^2}\delta_{ij} - \frac{\langle\bar{s}s\rangle}{12}\delta_{ij} - \frac{x^2}{192} \\
& \cdot m_0^2\langle\bar{s}s\rangle\left[1 - i\frac{m_s}{6}\not{x}\right]\delta_{ij} \\
& + \frac{i}{3}\left[\not{x}\left(\frac{m_s}{16}\langle\bar{s}s\rangle - \frac{1}{12}\left\langle u^\mu\Theta_{\mu\nu}^f u^\nu\right\rangle\right)\right. \\
& \left. + \frac{1}{3}(u\cdot x)\not{u}\left\langle u^\mu\Theta_{\mu\nu}^f u^\nu\right\rangle\right]\delta_{ij} - \frac{ig_sG_{ij}^{\alpha\beta}}{32\pi^2x^2}\left(\not{x}\sigma_{\alpha\beta}\right. \\
& \left. + \sigma_{\alpha\beta}\not{x}\right),
\end{aligned} \tag{15}$$

where $G_{ij}^{\alpha\beta} \equiv G_A^{\alpha\beta}t_{ij}^A$ is the external gluon field, $t_{ij}^A = \lambda_{ij}^A/2$ with λ_{ij}^A Gell-Mann matrices, $A = 1, 2, \dots, 8$ symbolizes color indices, m_s implies the strange quark mass, u_μ is the four-velocity of the heat bath, $\langle \bar{q}q\rangle$ is the temperature-dependent light quark condensate, and $\Theta^f_{\mu\nu}$ is the fermionic part of the energy-momentum tensor. Furthermore, the gluon condensate related to the gluonic part of the energy-momentum tensor $\Theta^g_{\alpha\beta}$ is expressed via relation [43]:

$$\left\langle Tr^c G_{\alpha\beta} G_{\lambda\sigma}\right\rangle_T = \left(g_{\alpha\lambda} g_{\beta\sigma} - g_{\alpha\sigma} g_{\beta\lambda}\right) A - \left(u_\alpha u_\lambda g_{\beta\sigma} - u_\alpha u_\sigma g_{\beta\lambda} - u_\beta u_\lambda g_{\alpha\sigma} + u_\beta u_\sigma g_{\alpha\lambda}\right) B, \tag{16}$$

where A and B coefficients are

$$A = \frac{1}{24}\left\langle G^a_{\alpha\beta} G^{a\alpha\beta}\right\rangle_T + \frac{1}{6}\left\langle u^\alpha \Theta^g_{\alpha\beta} u^\beta\right\rangle_T, \quad B = \frac{1}{3}\left\langle u^\alpha \Theta^g_{\alpha\beta} u^\beta\right\rangle_T. \tag{17}$$

In order to remove contributions originating from higher states, the standard Borel transformation with respect to q_0^2 is applied in the QCD side as well. By equating the coefficients of the selected structure $g_{\mu\nu}$ in both physical and QCD sides, and by employing the quark hadron duality ansatz up to a temperature-dependent continuum threshold $s_0(T)$, the final sum rules for Y_b are derived as

$$m^2_{Y_b}(T)\,\lambda^2_{Y_b}(T)\,e^{-m^2_{Y_b}(T)/M^2} = \int_{4(m_b+m_s)^2}^{s_0(T)} ds\,\rho^{\rm QCD}(s,T)\,e^{-s/M^2}. \tag{18}$$

To find the mass via TQCDSR, one should expel the hadronic coupling constant from the sum rules. It is commonly done by dividing the derivative of the sum rule given in (18) with respect to $(-M^{-2})$ to itself. Following these steps, the temperature-dependent mass is obtained as

$$m^2_{Y_b}(T) = \frac{\int_{4(m_b+m_s)^2}^{s_0(T)} ds\, s\rho^{\rm QCD}(s,T)\,e^{-s/M^2}}{\int_{4(m_b+m_s)^2}^{s_0(T)} ds\,\rho^{\rm QCD}(s,T)\,e^{-s/M^2}}, \tag{19}$$

where the thermal continuum threshold $s_0(T)$ is related to continuum threshold s_0 at vacuum via relation

$$s_0(T) = s_0\left[1 - \left(\frac{T}{T_c}\right)^8\right] + 4\left(m_b+m_s\right)^2\left(\frac{T}{T_c}\right)^8 \tag{20}$$

[44, 45]. For compactness, the explicit forms of spectral densities are presented in Appendix.

3. Phenomenological Analysis

In this section, the phenomenological analysis of sum rules obtained in (18) and (19) is presented. First, the input parameters and temperature dependance of relevant condensates are given. Following, the working regions of the obtained sum rules at vacuum are analyzed. The behavior of QCD sum rules at $T = 0$ is used to test the reliability of our analysis.

3.1. Input Parameters. During the calculations, input parameters given in Table 1 are used. In addition to these input parameters, temperature-dependent quark and gluon condensates, and the energy density expressions are necessary. The thermal quark condensate is chosen as

$$\langle \bar{q}q\rangle = \frac{\langle 0|\bar{q}q|0\rangle}{1+\exp\left(18.10042\left(1.84692\left[1/\text{GeV}^2\right]T^2 + 4.99216\left[1/\text{GeV}\right]T - 1\right)\right)}, \tag{21}$$

where $\langle 0|\bar{q}q|0\rangle$ is the light quark condensate at vacuum and which is credible up to a critical temperature $T_c = 190$ MeV. The expression given in (21) is obtained in [46, 47] from the Lattice QCD results given in [48, 49]. The temperature-dependent gluon condensate is parameterized via [46, 50]

$$\left\langle G^2\right\rangle = \left\langle 0\left|G^2\right|0\right\rangle \cdot\left[1 - 1.65\left(\frac{T}{T_c}\right)^{8.735} + 0.04967\left(\frac{T}{T_c}\right)^{0.7211}\right], \tag{22}$$

where $\langle 0|G^2|0\rangle$ is the gluon condensate in vacuum state and $G^2 = G^A_{\alpha\beta}G^{\alpha\beta}_A$. Additionally, for the gluonic and fermionic parts of the energy density, the following parametrization is used [46]

$$\left\langle \Theta^g_{00}\right\rangle = \left\langle \Theta^f_{00}\right\rangle = T^4 \exp\left(113.867\left[\frac{1}{\text{GeV}^2}\right]T^2 - 12.190\left[\frac{1}{\text{GeV}}\right]T\right) - 10.141\left[\frac{1}{\text{GeV}}\right]T^5, \tag{23}$$

which is extracted from the Lattice QCD data in [51].

3.2. Analysis of Sum Rules at $T = 0$. In order to get reliable results, obtained sum rules should be tested at vacuum, and the working regions of the parameters s_0 and M^2 should be determined. Within the working regions of s_0 and M^2, convergence of OPE and dominance of pole contributions should be assured. In addition, the obtained physical results should be independent of small variations of these parameters. Convergence of the OPE is tested by the following

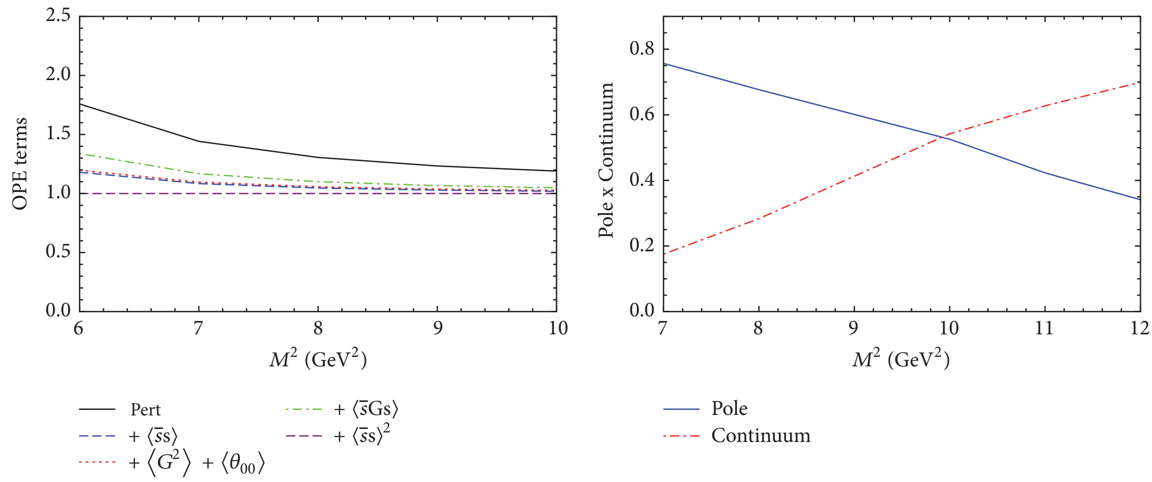

Figure 1: The OPE convergence of the sum rules: the ratio of the sum of the contributions up to specified dimension to the total contribution is plotted with respect to M^2 at $s_0 = 134$ GeV^2, $T = 0$ (left). Pole dominance of the sum rules: relative contributions of the pole (blue) and continuum (red-dashed) versus to the Borel parameter M^2 at $s_0 = 134$ GeV^2, $T = 0$ (right).

Table 1: Input parameters [20–23].

$m_s = (0.13 \pm 0.03)$ MeV
$m_b = (4.24 \pm 0.05)$ GeV
$m_0^2 = (0.8 \pm 0.2)$ GeV^2
$\langle s\bar{s}\rangle = -0.8 \times (0.24 \pm 0.01)^3$ GeV^3
$\left\langle 0 \left\| \frac{1}{\pi}\alpha_s G^2 \right\| 0 \right\rangle = (0.022)$ GeV^4

Table 2: Results obtained in this work for the mass of Y_b at $T = 0$, in comparison with literature.

	m_{Y_b} (MeV)
Present Work	10735^{+122}_{-107}
Experiment	$10889.9^{+3.2}_{-2.6}$ [20]
QCDSR	10880 ± 130 [11]
	10910 ± 70 [10]

criterion. The contribution of the highest order operator in the OPE should be very small compared to the total contribution. In Figure 1, the ratio of the sum of the terms up to the specified dimension to the total contribution is plotted to test the OPE convergence. It is seen that all higher order terms contribute less than the perturbative part for $M^2 \geq 6$ GeV^2. On the other hand, dominance of the pole contribution is tested as follows. The contribution coming from the pole of the ground state should be greater than the contribution of the continuum. In this work, the aforementioned ratio is

$$\mathrm{PC} = \frac{\Pi\left(M_{\max}^2, s_0\right)}{\Pi\left(M_{\max}^2, \infty\right)} > 0.50, \qquad (24)$$

when $M^2 \leq 10\,\mathrm{GeV}^2$ as can be seen in Figure 1. After checking these criteria, the working regions of the parameters M^2 and s_0 are determined as

$$6\ \mathrm{GeV}^2 \leq M^2 \leq 10\ \mathrm{GeV}^2; \qquad 132\ \mathrm{GeV}^2 \leq s_0 \leq 134\ \mathrm{GeV}^2, \qquad (25)$$

which is also consistent with $s_0 \simeq (m_H + 0.5\,\mathrm{GeV})^2$ norm [10]. Within these working regions, the variations of the mass of Y_b with respect to M^2 and s_0 are plotted in Figure 2. It is seen that the mass is stable with respect to variations of M^2 and s_0.

4. Results and Discussions

Following the analysis presented in previous section, the mass of the ground state estimated by the tetraquark current given in (3) at $T = 0$ is presented together with other results from literature in Table 2. It is seen that our results agree with other theoretical estimates and also with the experimental data on $Y_b(10890)$ [10, 11, 20]. Thus, the broad resonance Y_b can be described by the tetraquark current given in (3), and our analysis can be extended to finite temperatures.

To analyze the thermal properties of $Y_b(10890)$ resonance, the temperature dependencies of the mass and the pole residue of Y_b are plotted in Figure 3. It is seen that the mass and the pole residue of Y_b stay monotonous until $T \cong 0.12$ GeV. However, after this point, they begin to decrease promptly with increasing temperature. At the vicinity of the critical (or so called deconfinement) temperature, the mass reaches nearly 87% of its vacuum value. On the other hand, the pole residue decreased to 44% of its value at vacuum as shown in Figure 3.

Our predictions presented in Figure 3 are in good agreement with other QCD sum rules analysis on thermal behaviors of conventional or exotic hadrons [35–38]. However in [19], authors predicted a decrease of 5% in the mass of molecular 1^{++} state, and no change in the mass of charmonium 2^{-+} state, even beyond Hagedorn temperature $T_H \sim 177$ MeV. Since the decay properties also depend on temperature, and while the mass and pole residue diminish,

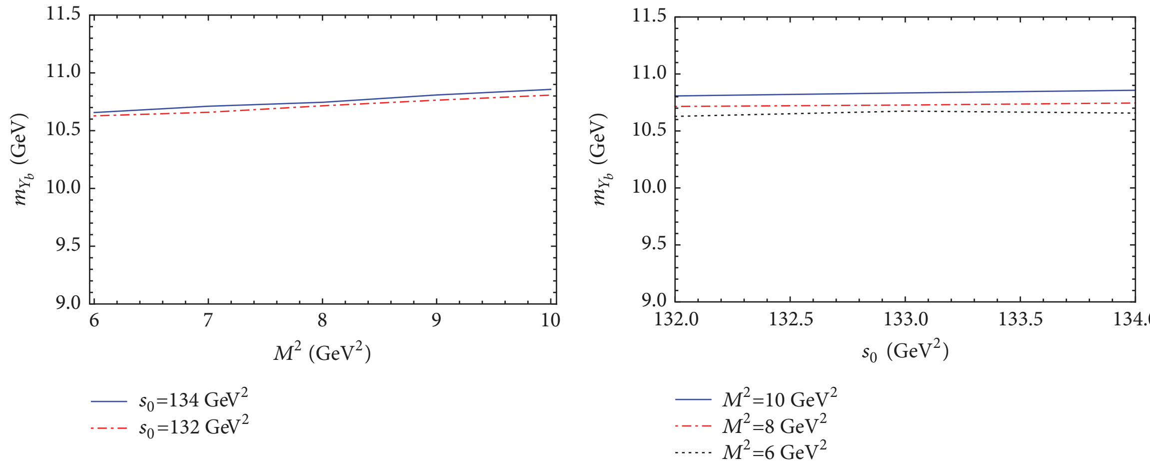

Figure 2: Mass of Y_b as a function of M^2 (left) and s_0 (right).

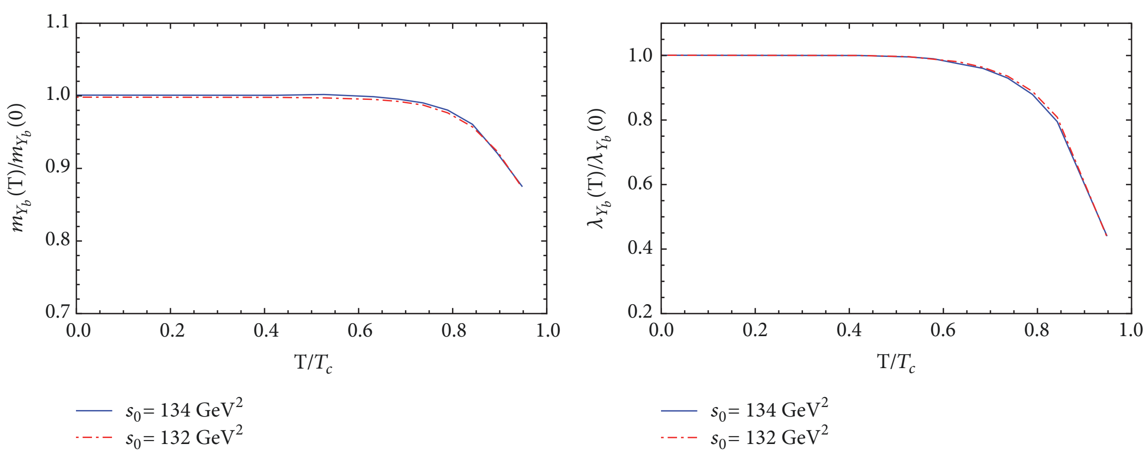

Figure 3: The mass (left) and pole residue (right) of Y_b as a function of temperature.

the decay width might increase with increasing temperature [52] and decay widths at finite temperature should also be investigated. However the current status of Y_b resonance is very complicated, since it is very close to $\Upsilon(5S)$ state. Thus studying its decays requires establishment of a good model in the hidden-beauty sector.

Finally, we would like to highlight the following remarks:

(i) We observed that the mass (the pole residue) of exotic $Y_b(10890)$ state starts to decrease near $T/T_c \simeq 0.7$.

(ii) Both quantities tend to diminish with increasing temperature up to critical temperature T_c.

(iii) Even though the sum rules at $T = 0$ estimates the mass of Y_b consistent with experimental data, more theoretical efforts are required to discriminate Y_b and $\Upsilon(5S)$.

(iv) In order to get more reliable results on tetraquarks from QCD sum rules, $O(\alpha_s^2)$ contributions suggested in [36] should be investigated.

In summary, we revisited the hidden-beauty exotic state Y_b and studied its properties at vacuum and finite temperatures. To describe the hot medium effects to the hadronic parameters of the resonance Y_b, TQCDSR method is used considering contributions of condensates up to dimension six. Our results for $T = 0$ are in reasonable agreement with the available experimental data and other QCD studies in the literature. Numerical findings show that Y_b can be well described by a scalar-vector tetraquark current. In the literature, remarkable drop in the values of the mass and the pole residue in hot medium was regarded as the signal of the QGP, which is called the fifth state of matter, phase transition. We hope that precise spectroscopic measurements in the exotic bottomonium sector can be done at Super-B factories, and this might provide conclusive answers on the nature and thermal behaviors of the exotic states.

Appendix

Thermal Spectral Density $\rho^{QCD}(s,T)$ for Y_b State

In this appendix, the explicit forms of the spectral densities obtained in this work are presented. The expressions for $\rho^{\text{pert.}}(s)$ and $\rho^{\text{nonpert.}}(s,T)$ areshown below as integrals over the Feynman parameters z and w, where θ is the step function,

$$\rho^{\text{pert.}}(s)=\frac{1}{3072\pi^6}\int_0^1 dz\int_0^{1-z}dw\frac{1}{\kappa^8\xi^2}\left\{\left[-\kappa m_b^2(z+w)+szw\xi\right]^2\left[\kappa^2 m_b^4 zw(z+w)\left[w^2+(-1+w)w+z(-1+4w)\right]\right.\right.$$
$$-2\kappa m_b^2\left[6m_s^2\Phi^2\left[7(-1+z)z(-7+8z)w+7w^2\right]+sz^2w^2\left(12(-1+z)z+(-12+25z)w+12w^2\right)\right]\xi \quad \text{(A.1)}$$
$$\left.\left.+szw\left(12m_s^2\Phi^2+35sz^2w^2\right)\xi^3\right]\right\}\theta\left[L(s,w,z)\right],$$

$$\rho^{\langle s\bar{s}\rangle}(s,T)=\frac{\langle\bar{s}s\rangle}{128\pi^4}\int_0^1 dz\int_0^{1-z}dw\frac{1}{\kappa^6}\left\{\left[\kappa^3 m_b^5 w(z+w)^2+\kappa^2 m_b^4 m_s(z+w)\left[19z^4+19(-1+w)^2w^2+2z(-1+w)w(-19+25w)\right.\right.\right.$$
$$\left.+z^3(-38+50w)+z^2\left[19+w(-88+81w)\right]\right]-\kappa^2 m_b^3 z(z+w)\left(m_s^2\Phi+2sw^2\right)\xi$$
$$-\kappa m_b^2 m_s szw\left[22z^4+22(-1+w)^2w^2+z^3(-44+111w)+z(-1+w)w(-44+111w)\right. \quad \text{(A.2)}$$
$$\left.+z^2\left[22+w(-155+197w)\right]\right]\xi+\kappa m_b sz^2w\left(m_s^2\Phi+sw^2\right)\xi^2$$
$$\left.\left.+3m_s s^2z^2w^2\left(z^2+(-1+w)w+z(-1+21w)\right)\xi^3\right]\right\}\theta\left[L(s,w,z)\right],$$

$$\rho^{\langle G^2\rangle+\langle\Theta_{00}\rangle}(s,T)=\frac{1}{4608\pi^2}\int_0^1 dz\int_0^{1-z}dw\frac{1}{\kappa^6\xi^2}\left\{192\pi^2\left\langle\Theta_{00}^f\right\rangle zw\xi^2\left[m_b^4\Phi^2(z+w)\right.\right.$$
$$\left.\times\left[3(-1+z)z+(-3+5z)w+3w^2\right]+m_b^2\Phi szw\left[-27(-1+z)z+27w-53zw-27w^2\right]\xi+30s^2z^2w^2\xi^3\right]$$
$$-g_s^2\left\langle\Theta_{00}^g\right\rangle\left[3m_b^4\Phi^2zw(z+w)\left[2(-1+z)^2z^2+(-1+z)z(-4+3z)w+(-2+z)(-1+3z)w^2+(-4+3z)w^3+2w^4\right]\right.$$
$$-m_b^2\Phi sz^2w^2\left[24(-1+z)^2z^2+(-1+z)z(-48+85z)w+\left[24+z(-133+121z)\right]w^2+(-48+85z)w^3+24w^4\right]\xi$$
$$\left.-12m_b^3m_s\Phi^3(z+w)^3\xi^2+12m_bm_s\Phi^2szz\left(z^2+10zw+w^2\right)\xi^3+30s^2z^3w^3(z+w)\xi^4\right]$$
$$+\left\langle\frac{\alpha_sG^2}{\pi}\right\rangle\pi^2\left[\kappa m_b^2\left\{-6m_s^2\Phi^2\left[5(-1+z)z^3+z^3w+(-5+z)w^3+5w^4\right]\right.\right. \quad \text{(A.3)}$$
$$\left.+sz^2w^2\left[2(-1+z)z^2(-18+11z)+z\left[72+z(-221+133z)\right]w+\left[36+z(-221+231z)\right]w^2(-58+133z)w^3+22w^4\right]\right\}\xi$$
$$-36\kappa^3m_b^3m_s(z+w)\left(z^2-6zw+w^2\right)\xi^2-60s^2z^3w^3(z+w)\xi^4$$
$$+\xi^2m_bzw\left[m_b^3(z+w)\left(4(-1+z)z^2(-3+4z)+6z\left[4+z(-11+8z)\right]w+\left[12+11z(-6+5z)\right]w^24(-7+12z)^3+16w^4\right)\right.$$
$$\left.\left.\left.+12m_ss\left(3z^2-26zw+3w^2\right)\xi^3\right]\right]\right\}\theta\left[L(s,w,z)\right]$$

$$\rho^{\langle\bar{s}Gs\rangle}(s,T)=\frac{m_sm_0^2\langle\bar{s}s\rangle}{64\pi^4}\left\{\int_0^1 dz\left\{3m_b^2+s(z-1)z\theta\left[L'(s,z)\right]\right\}+\int_0^1 dz\int_0^{1-z}dw\frac{1}{3\kappa^5}\right. \quad \text{(A.4)}$$
$$\left.\times\left\{zw\xi\left[\kappa m_b^2\left[5z^2+5(w-1)w+z(11w-5)\right]-16szw\xi^2\right]\theta\left[L(s,w,z)\right]\right\}\right\},$$

$$\rho^{\langle\bar{s}s\rangle^2}(s,T)=\frac{\langle\bar{s}s\rangle^2}{5184\pi^4}\left\{\int_0^1 dz\left[648m_b^2\pi^2+g_s^2m_bm_sz+54\pi^2\left(5m_s^2+4s\right)(z-1)z\right]\theta\left[L'(s,z)\right]+g_s^2\int_0^1 dz\int_0^{1-z}\frac{dwzw\xi}{\kappa^5}\right. \quad \text{(A.5)}$$
$$\left.\times\left[3\kappa m_b^2\left[7z^2+7(w-1)w+z(-7+15w)\right]-64szw\xi^2\right]\theta\left[L(s,w,z)\right]\right\},$$

where explicit expressions of the functions $L(s,w,z)$ and $L'(s,z)$ are

$$L\left[(s,w,z)\right.=\kappa^{-2}(-1+w)\left[(-1+w)w^2\right.$$
$$\left.+2(-1+w)wz+(-1+2w)z^2-swz\xi+z^3m_b^2\right], \quad \text{(A.6)}$$

$$L'(s,z)=sz(1-z)-m_b^2. \quad \text{(A.7)}$$

The below definitions are used for simplicity:

$$\kappa=z^2+z(w-1)+(w-1)w,$$
$$\Phi=(z-1)w+(z-1)w+w^2, \quad \text{(A.8)}$$
$$\xi=z+w-1.$$

Data Availability

No data were used to support this study.

Conflicts of Interest

The authors declare that they have no conflicts of interest.

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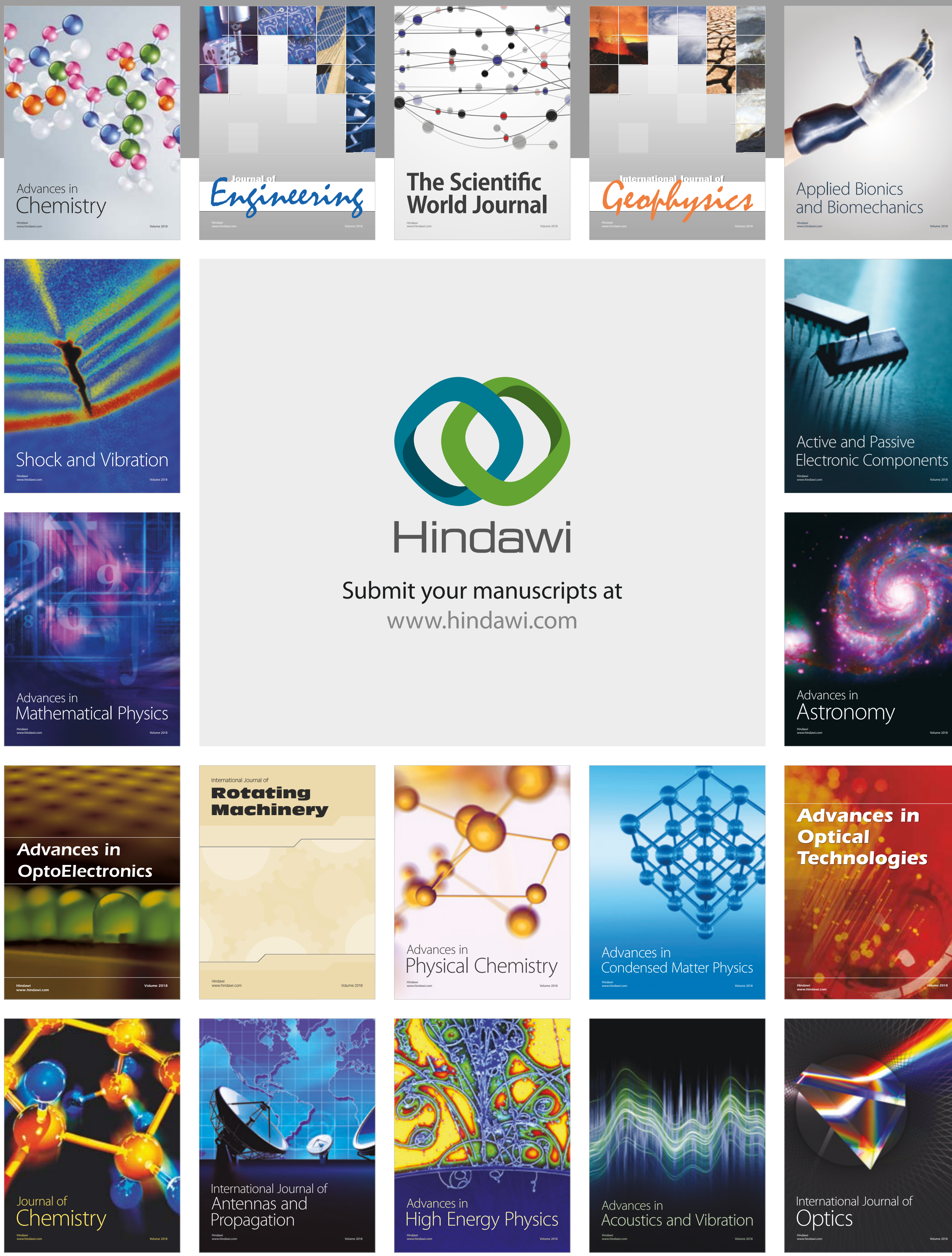
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