

The Scalar Hexaquark $uuddss$: a Candidate to Dark Matter?

K. Azizi,^{1,2,*} S. S. Agaev,³ and H. Sundu⁴

¹*Department of Physics, University of Tehran, North Karegar Ave., Tehran 14395-547, Iran*

²*Department of Physics, Doğuş University, Acibadem-Kadiköy, 34722 Istanbul, Turkey*

³*Institute for Physical Problems, Baku State University, Az-1148 Baku, Azerbaijan*

⁴*Department of Physics, Kocaeli University, 41380 Izmit, Turkey*

It is conventionally argued that Dark Matter (DM) has a non-baryonic nature, but if we assume that DM was frozen out before primordial nucleosynthesis and could not significantly impact primordial abundances this argument may be evaded. Then a hypothetical $SU(3)$ flavor-singlet, highly symmetric, deeply bound neutral scalar hexaquark $S = uuddss$, which due to its features has escaped from experimental detection so far, may be considered as a candidate for a baryonic DM. In the present work we calculate the mass and coupling constant of the scalar six-quark particle S by means of the QCD sum rule method. Our predictions for its mass are $m_S = 1180_{-26}^{+40}$ MeV ($m_s = 95$ MeV) and $\tilde{m}_S = 1239_{-28}^{+42}$ MeV ($m_s = 128$ MeV). Although these values of mass would produce thermally the cosmological DM abundance, existence of this state may contradict to stability of the oxygen nuclei, which requires further thorough analysis.

Introduction: From first days of the quark-parton model and Quantum Chromodynamics (QCD), hadrons with unusual quantum numbers and/or multiquark contents attracted interest of physicists. The conventional hadrons have quark-antiquark or three-quark compositions. Their masses and quantum parameters J^{PC} are in accord with predictions of this scheme and can be calculated using standard methods of particle physics. Unusual or exotic hadrons is expected to be built of four and more valence quarks or contain valence gluons. A main reason triggered intensive investigations of four-quark states was a mass hierarchy inside the lowest scalar multiplet, which found its explanation in the context of the four-quark model suggested by R. Jaffe [1]. Starting from 2003, i.e. from first observation of the exotic meson $X(3872)$ theoretical and experimental investigations of tetra and pentaquarks became one of the interesting and rapidly growing branches of high energy physics. Now, valuable experimental information collected during past years, as well as theoretical progress achieved to date, form two essential components of the exotic hadrons physics [2–6].

Another interesting result about multiquark hadrons with far-reaching consequences was obtained also by R. Jaffe [7]. He considered six-quark (dibaryon) states built of only light u , d , and s quarks that belong to flavor group $SU_f(3)$. By combining the color and spin of quarks and forming $SU_{cs}(6)$ "colorspin" group, Jaffe analyzed its representations and found that dibaryons only from the singlet and octet representations of $SU_f(3)$ may be light enough to be bound or resonant. Among six-quark states from these two representations $SU_f(3)$ singlet particles have zero spins. At the same time, all of singlet and octet dibaryons are strange particles, therefore structures containing merely u and d quarks cannot be bound and stable. Using the MIT quark-bag model

for analysis, Jaffe predicted existence of a H-dibaryon, i.e., of flavor-singlet and neutral six-quark $uuddss$ bound state with isospin-spin-parities $I(J^P) = 0(0^+)$. This double-strange six-quark structure with mass 2150 MeV lies 80 MeV below the $2m_\Lambda = 2230$ MeV threshold and is stable against strong decays. It can decay through weak interactions, which means that mean lifetime of H-dibaryon, $\tau \approx 10^{-10}$ s, is considerably longer than that of conventional mesons. The H-dibaryon with the mass obeying this limit is loosely-bound state, a subject to weak transformations.

The original work [7] was followed by numerous theoretical investigations, in which various models and methods of particle physics were used to calculate the H-dibaryon's mass [8–16]. As usual, results of these studies are controversial: thus, calculations in the framework of the corrected MIT bag model led to $m_H = 2240$ MeV which is just above the $2m_\Lambda$ threshold [8], whereas in a chiral model the authors [9] found $m_H = 1130$ MeV. To analyze $\Lambda - \Lambda$ interaction and estimate $\Lambda\Lambda$ binding energy other quark models were invoked as well [10, 12–14]. The H-dibaryon's mass extracted from the QCD two-point sum rules is consistent with the original result of Jaffe [15, 16]. In fact, m_H from Ref. [15] varies in limits 2.0 – 2.4 GeV and within an accuracy of the sum rule method $\sim 20\%$ agrees with the result of the quark-bag model. Calculations in Ref. [16] also confirmed existence of a bound state lying 40 MeV below the $2m_\Lambda$ threshold. The lattice simulations performed in Ref. [17] led to conclusion that m_H was below the $2m_N$ threshold 1880 MeV. In this paper the authors took into account the stability conditions of the nucleus and extracted $m_H \approx 1850$ MeV. The later lattice studies confirmed existence of a bound-state H-dibaryon, and predicted its binding energy ≈ 74.6 MeV [18] and (19 ± 10) MeV [19], respectively. In the context of the holographic QCD H-dibaryon was explored in Ref. [20], in which its mass was estimated about $m_H = 1.7$ GeV.

The hexaquark S (except for original papers, hereafter we use a hexaquark instead of a six-quark state, and de-

*Corresponding author

note it by S) was searched for by KTeV, Belle, and BaBar collaborations in exclusive $S \rightarrow \Lambda p \pi^-$ and inclusive $\Upsilon(1S)$ and $\Upsilon(2S)$ decays, in processes $\Upsilon(2S, 3S) \rightarrow S \overline{\Lambda \Lambda}$ [21–24]. All these experiments could not find evidence for the hexaquark S near the threshold $2m_\Lambda$ and were able only to impose limits on its mass m_S the latest being $m_S < 2.05$ GeV.

Recent activities around the S is inspired by renewed suggestions to consider it as a possible candidate to dark matter [25–29]. In accordance with this scenario if $m_S \leq 2(m_p + m_e) = 1877.6$ MeV the hexaquark is absolutely stable, because all possible decay channels of a free S is kinematically forbidden. For the m_S obeying the inequality $m_S < m_\Lambda + m_p + m_e = 2054.5$ MeV the hexaquark decays through a double-weak interaction, but even in this case its lifetime could be comparable with the age of the Universe. The lower bound of m_S is determined by a stability of ordinary nuclei, which are stable if $m_S > m_p + m_n + m_e - 2E$, where $2E$ is a binding energy of $p + n$. Then, for masses $1860 < m_S < 1880$ MeV, which assures a stability of the hexaquark and conventional nuclei, the S can explain both the relic abundance of the DM in the Universe and observed DM to ordinary matter ratio with less than 15% uncertainty [29]. But even S with the mass in the range $1.3 \lesssim m_S \lesssim 2m_p$ and with radius $(1/6 - 1/4)r_p$, where $r_p \approx 0.86$ fm is the a proton radius, is consistent with the stability of nuclei, with Λ decays in double-strange hypernuclei and experimental limits on existence of exotic isotopes of helium and other nuclei [25]. There are, however, objections to this picture connected with a production process of hexaquarks in the early-universe [30], or with observed supernova explosion [31].

The hexaquark S as a candidate to DM was recently analyzed in Ref. [32] as well. In this work the mass of S was evaluated by modeling it as a bound state of scalar diquarks. Using the effective Hamiltonian to describe dominant spin-spin interactions in diquarks [4], the authors expressed m_S in terms of constituent diquark masses m_{ij} and chromomagnetic couplings k_{ij} . The masses of diquarks and chromomagnetic couplings may be extracted from analysis of baryon spectroscopy. Alternatively, k_{ij} can also be fixed to reproduce masses of the light scalar mesons $f_0(500)$, $K^*(800)$, $f_0(980)$, and $a_0(980)$ interpreted as tetraquarks [33]. It turns out that spin-spin couplings in tetraquarks are about a factor of four larger compared to the spin-spin couplings in the baryons. Because the hexaquark itself is an exotic six-quark meson for calculation of m_S it is reasonable to employ parameters estimated from analysis of the light tetraquarks. Calculations carried out in Ref. [32] predict $m_S \approx 1.2$ GeV which reproduces the cosmological DM abundance, but may contradict to stability of oxygen nuclei.

Calculations: In the present work we calculate the mass of S by treating it as a bound state of three scalar diquarks. To this end, we employ the QCD two-point sum rule approach, which is one of the powerful nonper-

turbative methods to explore hadrons. As starting point, the method uses the correlation function

$$\Pi(p) = i \int d^4x e^{ipx} \langle 0 | \mathcal{T} \{ J(x) J^\dagger(0) \} | 0 \rangle, \quad (1)$$

and extract from its analysis sum rules to compute spectroscopic parameters of the hexaquark. The main ingredient of this analysis is the interpolating current $J(x)$ which we choose it in the following form

$$J(x) = \epsilon^{abc} [u^T(x) C \gamma_5 d(x)]^a [u^T(x) C \gamma_5 s(x)]^b \times [d^T(x) C \gamma_5 s(x)]^c, \quad (2)$$

where $[q^T C \gamma_5 q']^a = \epsilon^{amn} [q_m^T C \gamma_5 q_n']$ and a, b, c, m, n are color indices with C being the charge-conjugation operator.

As is seen, the hexaquark is composed of the scalar diquarks $[q^T C \gamma_5 q']^a$ in the color antitriplet and flavor antisymmetric states. These diquarks are most attractive ones [34], and six-quark mesons composed of them should be lighter and more stable than bound states of other two-quarks. Mathematical manipulations to derive sum rules for the mass and coupling of the hexaquark are carried out in accordance with standard prescriptions of the method. Thus, first we express the correlation function $\Pi(p)$ in terms of the hexaquark's mass m_S and coupling f_S , as well as its matrix element

$$\langle 0 | J | S \rangle = m_S f_S. \quad (3)$$

Separating from each another the ground-state term and contributions due to higher resonances and continuum states for $\Pi^{\text{Phys}}(p)$ we get

$$\Pi^{\text{Phys}}(p) = \frac{\langle 0 | J | S(p) \rangle \langle S(p) | J^\dagger | 0 \rangle}{m_S^2 - p^2} + \dots \quad (4)$$

The expression of the matrix element (3) allows us to rewrite $\Pi^{\text{Phys}}(p)$ in the form

$$\Pi^{\text{Phys}}(p) = \frac{m_S^2 f_S^2}{m_S^2 - p^2} + \dots, \quad (5)$$

where dots denote contributions of higher resonances and continuum states.

To calculate the QCD or OPE side of the sum rules, we insert the current $J(x)$ to Eq. (1), contract relevant quark fields and obtain $\Pi^{\text{OPE}}(p)$ in terms of the quark propagators:

$$\begin{aligned} \Pi^{\text{OPE}}(p) = & i \delta^{af} \delta^{a'f'} \delta^{bd} \delta^{b'd'} \delta^{ce} \delta^{c'e'} \int d^4x e^{ipx} \{ \text{Tr} [S_d^{ee'}(x)] \\ & \times \gamma_5 \tilde{S}_s^{ff'}(x) \gamma_5 \} \text{Tr} [S_u^{aa'}(x) \gamma_5 \tilde{S}_d^{bb'}(x) \gamma_5] \\ & \times \text{Tr} [S_u^{cc'}(x) \gamma_5 \tilde{S}_s^{dd'}(x) \gamma_5] \} + 511 \text{ similar terms}, \end{aligned} \quad (6)$$

where $\tilde{S}(x) = CS^T(x)C$. To proceed, we employ the x -space light-quark propagator

$$\begin{aligned}
S_q^{ab}(x) = & i\frac{\not{x}}{2\pi^2 x^4} \delta_{ab} - \frac{m_q}{4\pi^2 x^2} \delta_{ab} - \frac{\langle \bar{q}q \rangle}{12} \left(1 - i\frac{m_q}{4} \not{x}\right) \delta_{ab} \\
& - \frac{x^2}{192} \langle \bar{q}g_s \sigma Gq \rangle \left(1 - i\frac{m_q}{6} \not{x}\right) \delta_{ab} \\
& - \frac{ig_s G_{ab}^{\mu\nu}}{32\pi^2 x^2} [\not{x}\sigma_{\mu\nu} + \sigma_{\mu\nu}\not{x}] - \frac{\not{x}x^2 g_s^2}{7776} \langle \bar{q}q \rangle^2 \delta_{ab} \\
& - \frac{x^4 \langle \bar{q}q \rangle \langle g_s^2 G^2 \rangle}{27648} \delta_{ab} + \frac{m_q g_s}{32\pi^2} G_{ab}^{\mu\nu} \sigma_{\mu\nu} \left[\ln\left(\frac{-x^2 \Lambda^2}{4}\right) + 2\gamma_E \right] \\
& + \dots, \tag{7}
\end{aligned}$$

where $q = u, d$ or s , $\gamma_E \simeq 0.577$ is the Euler constant, and Λ is a scale parameter. We also use the notations $G_{ab}^{\mu\nu} \equiv G_A^{\mu\nu} t_{ab}^A$, $A = 1, 2, \dots, 8$, and $t^A = \lambda^A/2$, with λ^A being the Gell-Mann matrices.

After inserting the light-quark propagators into Eq. (6), we get the correlation function $\Pi^{\text{OPE}}(p)$ in terms of QCD degrees of freedom. The next step is to perform the resultant Fourier integrals over four- x . Afterwards we equate the invariant amplitudes $\Pi^{\text{Phys}}(p^2)$ and $\Pi^{\text{OPE}}(p^2)$ to find the desired sum rule in momentum space. We apply the Borel transformation to both sides of the obtained sum rule to suppress contributions of the higher resonances and continuum states, and using the quark-hadron duality assumption, which is a quintessence of the sum rule method, perform the continuum subtraction. An equality derived after these manipulations, contains the mass and coupling constant of the particle S. To find the sum rules for m_S and f_S we need an extra expression which can be obtained by acting $d/d(-1/M^2)$ to the first equality. The sum rules for m_S and f_S derived by this way have perturbative and nonperturbative components. The latter contains vacuum condensates of various local quark, gluon, and mixed operators, which appear after sandwiching relevant terms in $\Pi^{\text{OPE}}(p)$ between vacuum states.

As the hexaquark is composed of six quarks, relevant computations are lengthy and time consuming. In Appendix we explain some details of calculations, and write down explicitly the Borel transformed and subtracted invariant amplitude $\Pi^{\text{OPE}}(M^2, s_0)$ including nonperturbative terms up to dimension ten. The full expression of $\Pi^{\text{OPE}}(M^2, s_0)$ contains terms up to dimension thirty, therefore we refrain from providing them here. In numerical computations, we take into account all these higher dimensional terms bearing in mind that they appear due to the factorization hypothesis as product of basic condensates.

In analyses and computations, we utilize the quark, gluon and mixed condensates

$$\begin{aligned}
\langle \bar{q}q \rangle &= -(0.24 \pm 0.01)^3 \text{ GeV}^3, \quad \langle \bar{s}s \rangle = 0.8 \langle \bar{q}q \rangle, \\
\langle \bar{q}g_s \sigma Gq \rangle &= m_0^2 \langle \bar{q}q \rangle, \quad \langle \bar{s}g_s \sigma Gs \rangle = m_0^2 \langle \bar{s}s \rangle, \\
m_0^2 &= (0.8 \pm 0.1) \text{ GeV}^2 \\
\langle \frac{\alpha_s G^2}{\pi} \rangle &= (0.012 \pm 0.004) \text{ GeV}^4, \tag{8}
\end{aligned}$$

which are determined at the scale $\mu = 1$ GeV. We work in the approximation $m_u = m_d = 0$, but keep a dependence on m_s . The parameter Λ is varied within the limits (0.5, 1) GeV.

Since the S is tightly bound state with the radius much smaller than usual hadrons, its parameters should be explored at relatively large momentum scales μ . Therefore, we calculate the mass and coupling of the hexaquark at $\mu = 2$ GeV which corresponds to the radius 0.1 fm. To reveal a sensitivity of m_S and f_S on the scale μ , we evaluate the same parameters at $\mu = 1$ GeV as well.

The mass of the strange quark $m_s = 95_{-3}^{+9}$ MeV in the \overline{MS} scheme and at the scale $\mu = 2$ GeV can be found in Ref. [35]. We evolve the condensates (8) to this scale and perform numerical computations. At the scale $\mu = 1$ GeV calculations are carried out by employing (8) and the mass m_s at this scale that differs from PDG value by a factor 1.35.

Another important problem is a proper choice for the the Borel M^2 and continuum threshold s_0 parameters. First of them has been introduced upon Borel transformation, the second one is necessary to separate the ground-state and continuum contributions from each another, as we previously have mentioned. These parameters are not arbitrary, but should meet the well-known requirements. Thus, at maximum value of the Borel parameter the pole contribution (PC) should constitute a fixed part of the correlation function, whereas at minimum of M^2 it must be a dominant contribution. We define PC in the form

$$\text{PC} = \frac{\Pi^{\text{OPE}}(M^2, s_0)}{\Pi^{\text{OPE}}(M^2, \infty)}, \tag{9}$$

and at M_{max}^2 impose on it a restriction $\text{PC} > 0.2$, which is usual for multi-quark hadrons. The minimum of the Borel parameter M^2 is fixed from convergence of the sum rules, i.e. at M_{min}^2 contribution of the last term (or a sum of last few terms) cannot exceed, for example, 0.01 part of the whole result. There is an another restriction on the lower limit M_{min}^2 : at this M^2 the perturbative contribution has to prevail over the nonperturbative one.

The sum rule predictions, in general, should not depend on the parameter M^2 . But in real calculations m_S and f_S demonstrate sensitiveness to the choice of M^2 and one should find a plateau where this dependence is minimal. The continuum threshold parameter s_0 separates a ground-state contribution from the ones arising from higher resonances and continuum states. Stated differently, s_0 should be below the first excited state of the hexaquark S. Parameters of conventional hadrons' excited states are known either from experimental measurements or from alternative theoretical studies. In the lack of similar information for multi-quark hadrons, one fixes s_0 to achieve a maximum for PC ensuring, at the same time, fulfilments of other constraints, and keeping under control a self-consistency of computations. The self-consistent analysis implies that a gap between the

mass m_S and the parameter $\sqrt{s_0}$ used for its extraction should be within reasonable limits of a few hundred MeV.

Performed analysis allows us to determine the working regions

$$M^2 \in [1.3, 1.6] \text{ GeV}^2, \quad s_0 \in [2.5, 2.9] \text{ GeV}^2, \quad (10)$$

which obey all aforementioned restrictions.

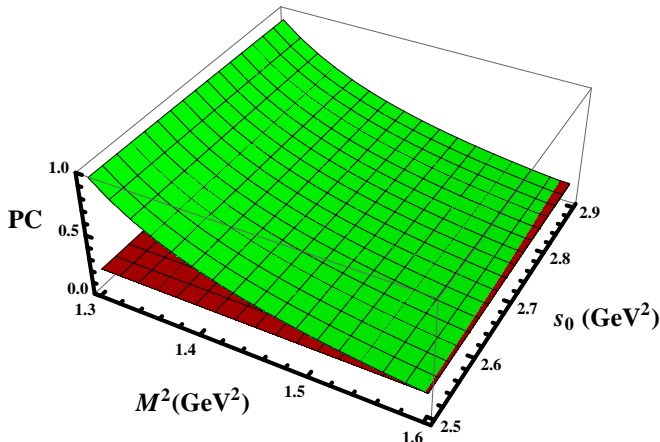


FIG. 1: Dependence of the pole contribution on M^2 and s_0 . The (red) surface $PC = 0.2$ is also shown.

In Fig. 1 we plot the pole contribution, when M^2 and s_0 are varying within limits (10): at $M^2 = 1.3$ the pole contribution is 0.9, whereas at $M^2 = 1.6$ it becomes equal to 0.2. The predictions for the mass m_S is pictured in Fig. 2, where a mild dependence on the parameters M^2 and s_0 is seen. The results for the spectroscopic parameters of the hexaquark S read: at the scale $\mu = 2$ GeV (for $m_s = 95$ MeV)

$$m_S = 1180_{-26}^{+40} \text{ MeV}, \quad f_S = 8.56_{-0.26}^{+0.03} \times 10^{-6} \text{ GeV}^7, \quad (11)$$

and at the scale $\mu = 1$ GeV (for $m_s = 128$ MeV)

$$\tilde{m}_S = 1239_{-28}^{+42} \text{ MeV}, \quad \tilde{f}_S = 9.18_{-0.22}^{+0.03} \times 10^{-6} \text{ GeV}^7. \quad (12)$$

Let us note that in computation of \tilde{m}_S and \tilde{f}_S the Borel parameter has been varied within the limits $M^2 \in [1.34, 1.63] \text{ GeV}^2$.

Theoretical uncertainties in the sum rule predictions (11) and (12) appear mainly due to working windows for the auxiliary parameters M^2 and s_0 , and the scale Λ . The ambiguities connected with various vacuum condensates are numerically small. We do not include into errors of the mass and coupling corrections generated by different choices of the scale μ , but keep (m_S, f_S) and $(\tilde{m}_S, \tilde{f}_S)$ as two sets of parameters. Let us note that variation of the mass $\Delta m_S(\mu) \approx 60$ MeV is comparable with other errors and does not exceed a few percent of m_S . It is not difficult to check also self-consistent character of obtained results. Indeed, estimating $\sqrt{s_0} - m_S$ we

get [400, 525] MeV, which can be accepted as a normal value for the mass difference between the ground-state and first excited hexaquarks.

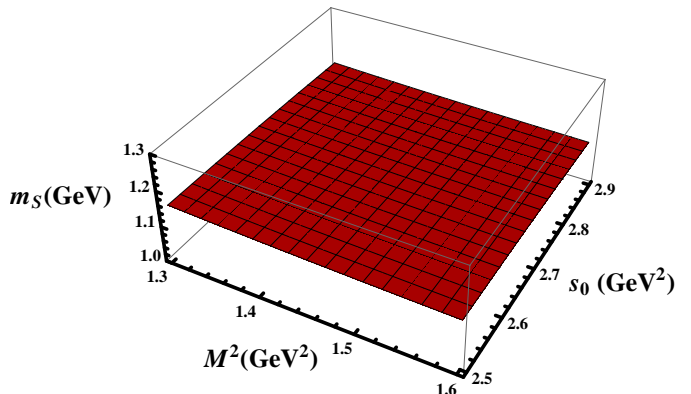


FIG. 2: The mass of the hexaquark S as a function of the Borel and continuum threshold parameters. The mass of the strange quark is $m_s = 95$ MeV.

Discussion and Conclusions: We have considered the spin-0, parity-even, highly symmetric S-hexaquark of $uuddss$ with $Q = 0$, $B = 2$ and $S = -2$. Using the technique of QCD sum rule, we have found that for the chosen interpolating current the intervals $M^2 \in [1.3, 1.6] \text{ GeV}^2$, $s_0 \in [2.5, 2.9] \text{ GeV}^2$ for the auxiliary parameters fulfill the requirements of the method discussed above. For these intervals, we could able to reach [0.9 – 0.2] pole contributions to the sum rules, and have extracted $m_S = 1180_{-26}^{+40}$ MeV ($m_s = 95$ MeV) and $\tilde{m}_S = 1239_{-28}^{+42}$ MeV ($m_s = 128$ MeV) for the mass of the S-hexaquark. This range of the mass implies that the hexaquark S is an absolutely stable particle.

It is worth noting that in the context of the sum rule method the scalar six-quark particle was investigated in Refs. [15, 16]. In Ref. [15], the authors explored different interpolating currents and carried out calculations by taking into account nonperturbative terms up to $\langle \bar{q}q \rangle^4$ and $m_s \langle \bar{q}q \rangle^5$ orders. Using $m_s = 0.2$ GeV, and an approximation $\langle \bar{q}q \rangle = \langle \bar{s}s \rangle$, for the mass of the hexaquark they found $m_S = 2.4$ GeV. But, because of uncertainties of calculations, the authors could not determine whether the mass of this particle lies above or below the $\Lambda\Lambda$ threshold, i.e. whether it is stable or not. In Ref. [16] the two-point correlation function was found by employing for the hexaquark a molecular type current. The authors took into account only terms proportional to $\langle \bar{q}q \rangle^2$ and $\langle \bar{q}q \rangle^4$, and neglected the gluon and mixed condensates. Prediction in Ref. [16] was made using the strange quark mass $m_s = 150$ MeV; it was found that the mass of the hexaquark is $m_S \approx 2.19$ GeV which corresponds to a bound state 40 MeV below the $\Lambda\Lambda$ threshold.

The accuracy of calculations performed in the present work considerably exceeds the accuracy of previous investigations. In this paper we have included into analysis not only the quark, but also the gluon and mixed

condensates. In computations we took into account non-perturbative terms up to dimension thirty, and for the mass of the strange quark used its contemporary value $m_s = 95_{-3}^{+9}$ MeV. Therefore, our result for the mass of S-hexaquark differs from estimations of Refs. [15, 16], but is in accord with output of the chiral model [9]. Our prediction for m_S (and \tilde{m}_S) almost coincides with one made recently in Ref. [32], in which it was obtained by modeling the S-hexaquark as a bound state of scalar diquarks and using recent progress in theoretical and experimental physics of multi-quark hadrons.

There are a lot of constraints on the mass and radius of the hexaquark as a candidate to the DM. They stem from analyses of the different production modes of S-hexaquark in high energy hadron and e^-e^+ collisions, from its interactions with ordinary baryons, and from analysis of cosmological processes [27, 28]. As we have noted above, accelerator experiments could not observe the hexaquark S near the threshold $2m_\Lambda$ and only imposed limits on its mass $m_S < 2.05$ GeV. There are some reasons that make problematic detection of S-hexaquark at colliders. In fact, the hexaquark is neutral and compact particle and its mass is close to the mass of neutron, which makes difficult to separate relevant signals from ones generated by neutron. Additionally, as a flavor-singlet particle, the S-hexaquark presumably does not couple to a hadronic content of photon, pion and other flavor-nonsinglet mesons, or such interactions are very weak. It is possible that due to these features it escapes detection in hadron collisions. Strategies for discovering a stable hexaquark in various hadronic processes and relevant problems were discussed in Ref. [27]. In accordance with estimates of this work, within existing experimental datasets should be a few hundred events with S or its antiparticle \bar{S} . For creation of the S-hexaquark e^-e^+ collisions may be more promising than hadronic processes, because flavor-singlet multi-gluon states copiously produced in these collisions eventually may transform to S particles [28]. In any case, the constraint $m_S < 2.05$ GeV extracted from collider processes does not contradict to existence of a stable particle with the mass $m_S \approx 1.2$ GeV.

As a particle carrying the baryon number $B = 2$, the S-hexaquark interacts with other baryons and nuclei, and these processes might be utilized to detect it. Parameters of such interactions have to be calculated using perturbative or nonperturbative methods of QCD, and confronted with available experimental data. There are some active experiments designed for direct detection of DM scattering on target materials of detectors [36–38]. Recently the XENON and SuperCDMS collaborations reported about constraints on light DM-nucleon scattering cross-section σ_{DMN} extracted from their studies [37, 38]. In accordance with Ref. [37], for DM particles with the mass ~ 1 GeV and a spin-independent $DM - N$ interaction the cross-section σ_{DMN} is $\sigma_{DMN} \approx 10^{-38}$ cm²; and $\sigma_{DMp} \approx 10^{-32}$ cm² and $\sigma_{DMn} \approx 10^{-33}$ cm² if this interaction is spin dependent (see, Fig. 5 in Ref. [37]). The

information provided by the SuperCDMS (see, Fig. 13 of [38]) allows us to estimate $\sigma_{DMN} \sim 10^{-36}$ cm². The limit on DM-proton cross-section $\sigma_{DMp} \leq 0.6 \times 10^{-30}$ cm² was extracted also in Ref. [39], in which the authors relied on XDC data [40]. Results of other experiments devoted to DM-nucleon scattering can be found in Refs. [37–39].

The theoretical investigation of the hexaquark-nucleon scattering cross-section σ_{NS} should be performed in the context of QCD, i.e., in a framework of the Standard Model (SM). In this aspect, this task differs considerably from the situation in theories, where DM particle and a mediator of DM-nucleon interaction are introduced at expense of various extensions of SM [41]. Because the hexaquark is the particle composed of the conventional quarks, a mediator of S – N interaction also should belong to ordinary hadron spectroscopy. It should be neutral flavor-singlet scalar particle with a mass < 1 GeV: There are few candidates to play a role of such mediator. The scalar mesons σ [in new classification $f_0(500)$], $f_0(980)$, and a scalar glueball G may couple to both S and baryons, and carry this interaction. The σ and $f_0(980)$ are singlets from the lowest nonet of scalar mesons. Difficulties in interpretation of these mesons as $q\bar{q}$ states inspired suggestion about their diquark-antidiquark nature [1]. Within this paradigm problems with low masses, and a mass hierarchy inside the light nonet seem found their solutions. The current status of relevant theoretical studies can be found in Refs. [42–44]. The scalar glueball G , in accordance with various estimations, has the mass $\gtrsim 1$ GeV, but due to mixing with quark component would appear as a part of σ and $f_0(980)$ mesons [45].

Here, for the sake of concreteness, we analyze only σ -exchange processes. Then, to evaluate the cross-section σ_{NS} one has to compute strong couplings corresponding to vertices $NN\sigma$ and $SS\sigma$: This is necessary to evaluate σ_{NS} using one of QCD approaches. The $NN\sigma$ coupling, actually, is known, and was calculated in the context of the sum rule method in Refs. [46, 47]. In these articles the meson σ was treated either as mixed $q\bar{q} + G$ or pure $q\bar{q}$ states. The coupling $NN\sigma$, where σ is the diquark-antidiquark exotic meson or admixture of four-quark and glueball components, as well as $SS\sigma$ were not calculated using available methods of QCD. It is clear, that the hexaquark S is self-interacting particle, which takes place through the same σ -exchange mechanism. In other words, the S-hexaquark as the DM particle belongs to class of self-interacting DM models. The cross-section σ_{SS} for elastic S – S interaction can be computed using the strong coupling of the vertex $SS\sigma$. It is worth noting that σ_{NS} was evaluated in Ref. [25] by modeling S – N interaction via one-boson exchange Yukawa potential. The result obtained there put on σ_{NS} the limit $\sigma_{NS} \lesssim 10^{-3}$ mb or $\lesssim 10^{-30}$ cm². At high velocities of the interacting particles $v \approx c$, naive estimates for σ_{NS} and σ_{SS} led to constraints $\sigma_{NS} \geq 0.25\sigma_{NN} \approx 5$ mb and $\sigma_{SS} \geq 0.25\sigma_{NS} \approx 1.25$ mb [28], respectively. The restrictions on σ_{DMp} for spin-dependent interactions 10^{-27} cm² $< \sigma_{DMp} < 10^{-24}$ cm² were obtained in Ref.

[48]. A strong theoretical analysis, in accordance with the scheme outlined above, is required to make a reliable prediction for σ_{NS} compatible or not with existing experimental limits.

Besides direct detection experiments, there are outer space cosmic ray (CR) experiments, results of which can be utilized to explore the Dark Matter through its decays to SM particles. Such investigations may be helpful also for studying matter-antimatter asymmetry in the Universe. The finished satellite-based PAMELA and ongoing AMS-02 experiments provided information useful for these purposes. The PAMELA was constructed to detect galactic CRs, and mainly their positron and antiproton components [49]. Measurements revealed a rise in the positron-electron ratio $e^+/(e^- + e^+)$ at energies above 10 GeV. The abundance of positrons in CRs was confirmed by AMS-02 up to energies 350 GeV [50] and later till 500 GeV. The similar enhancement was observed in antiproton to proton ratio \bar{p}/p , as well [51, 52]. A standard scenario implies that antiparticles are produced due to inelastic interactions of CR nuclei with particles of the interstellar gas. But rates of such processes are small, therefore deviation of these ratios from expected small values may be interpreted in favor of DM decaying to SM particles. Collected data on antiparticle anomaly in the Universe provide valuable information to verify different DM models, and impose constraints on DM-SM interactions [41]. Alternatively, the observed abundance of high-energy positrons in CRs may be connected with accelerating effects of nearby sources, such as a pulsar, supernova remnants [53–55]. This mechanism implies some anisotropy in detected antiparticle fluxes, whereas data are consistent with their isotropic distributions: this is significant obstacle in attempts to explain antiparticle excess in CRs using local sources.

The hexaquark S produces leptons in the $S\bar{S}$ annihilation and inelastic interactions with particles of the interstellar gas. Productions of various groups of σ , $f_0(980)$, π , and K mesons would be main channels at low energy annihilations. At high energies these light mesons should be accompanied by numerous baryons. Dominant decay modes of the mesons σ , $f_0(980)$, π , and K , and relevant branching ratios are well known [35], which can be employed to estimate production rates of electrons and positrons to account for discussed effects.

Important restrictions on the mass of the hexaquark arise from observed cosmological abundance of the DM, and DM to ordinary matter ratio in the Universe. The concept of DM is necessary (excluding models with modified theory of gravity) to account for observed astrophysical effects, such as rotation curves of galaxies, gravitational lensing, and other phenomena. Direct evidence for the existence of DM came from analysis of bullet clusters [56]. All these phenomena can be explained within the concept of DM provided DM particles are stable. The hexaquark with mass $m_S \approx 1.2$ GeV reproduces observed DM abundance and DM/matter ratio, while a larger m_S gives a smaller relic abundance [32]. The S-hexaquark

is the stable particle, and its self-interaction and elastic scattering on ordinary matter does not reduce the total mass of DM in some location, which is crucial to describe aforementioned phenomena. Of course, interaction of S with baryon and photon fluids may alter matter power spectrum and the cosmic microwave background, but they are not strong enough to generate significant effects [32]. Contrary, the S-baryon (i.e., DM-galactic gas) interactions may produce a DM disk embedded within the spherical galactic halo, lead to co-rotation of DM with the gas and forming DM density structure similar to that of the gas [28]. Parameters of the DM disk in a galaxy, its thickness, for example, depend on the mass of DM and the cross-section σ_{DMN} [57].

Another constraint on m_S is connected with stability of conventional nuclei. The reason is that very small mass of the hexaquark may contradict to stability of existing nuclei, because for small m_S nucleons inside nuclei would bind to the S state faster than what is allowed by observed stability of these nuclei. This process runs through double-weak production of the off-shell Λ^* baryons by a pair of nucleons pn , pp or nn . Because our estimate for the mass of the hexaquark is $m_S \approx 1.2$ GeV, the main sources of the virtual Λ^* baryons are the weak decays $p \rightarrow \Lambda^* \pi^+$, $n \rightarrow \Lambda^* \pi^0$, and an internal conversion (udd) \rightarrow (uds). Generated by this way virtual Λ^* s afterwards through the strong-interaction process $\Lambda^* \Lambda^* \rightarrow S$ form the hexaquark S.

The matrix element of the reaction $NN \rightarrow SX$ can be written as a product of the amplitude for the nucleons' double-weak transitions to a pair of virtual Λ^* s, and matrix element for creation of the S from the Λ^* s [26]

$$\mathcal{M}(NN \rightarrow SX) \approx \mathcal{M}(NN \rightarrow \Lambda^* \Lambda^* X) \mathcal{M}(\Lambda^* \Lambda^* \rightarrow S). \quad (13)$$

Then lifetime of the nucleus \mathcal{N} decaying to \mathcal{N}' and the hexaquark is

$$\tau(\mathcal{N} \rightarrow \mathcal{N}' SX) \simeq \frac{3\text{yr}}{|\mathcal{M}(\Lambda^* \Lambda^* \rightarrow S)|^2}, \quad (14)$$

which should be confronted with the Super-Kamiokande (SK) limit for the oxygen nuclei

$$\tau(^{16}\text{O}_8 \rightarrow \mathcal{N}' SX) \gtrsim 10^{26}\text{yr}. \quad (15)$$

Equation (14) is rather rough estimate for τ , which is seen from treatment, for instance, of the matrix element $|\mathcal{M}(\Lambda \rightarrow N)|^2$ used to derive it [26]). This matrix element was calculated there in the harmonic oscillator model, and was also inferred from phenomenological analysis: obtained results differ from each other by approximately 10 times. The prediction (14) was made by employing an average value of $|\mathcal{M}(\Lambda \rightarrow N)|^2$.

The situation with $|\mathcal{M}(\Lambda^* \Lambda^* \rightarrow S)|^2$ is even worst than in the previous case. The matrix element $|\mathcal{M}(\Lambda^* \Lambda^* \rightarrow S)|^2$ describes the strong process $\Lambda^* \Lambda^* \rightarrow S$ and plays a crucial role in theoretical estimations of the $^{16}\text{O}_8$ lifetime.

In Ref. [26] the $\mathcal{M}(\Lambda^*\Lambda^* \rightarrow S)$ was calculated as overlap integral of the final hexaquark and initial $\Lambda^*\Lambda^*$ baryons's wave functions, the latter being factored into wave functions of the two Λ^* baryons, and a wave function of two nucleons inside nucleus. The Λ^* baryon and hexaquark S wave functions were written down using the Isgur-Karl (IK) nonrelativistic harmonic oscillator quark model [58] and its generalization to six-quark system [26]. These functions depend on parameters $\alpha_{B(S)} = 1/\sqrt{\langle r_{B(S)}^2 \rangle}$, where $\langle r_B^2 \rangle$ and $\langle r_S^2 \rangle$ are mean charge radii of the Λ^* baryon and hexaquark S, respectively. The features of the two nucleons inside nucleus were modeled in Ref. [26] using Brueckner-Bethe-Goldstone (BBG) wave function. A wide class of the two-nucleon wave functions including the BBG, the Miller-Spencer wave functions, and ones extracted from results of Ref. [59], was employed in Ref. [32] to study the stability of the oxygen nucleus.

The IK wave functions used to model the Λ^* baryons and the hexaquark S suffer from serious drawbacks. Thus, the parameter $\alpha_B = 0.406 \text{ fm}^{-1}$ necessary to describe the mass splitting of lowest lying $\frac{1}{2}^+$ and $\frac{3}{2}^+$ baryons corresponds to radius of the proton 0.49 fm, whereas to reproduce the experimental value 0.86 fm one needs $\alpha_B = 0.221 \text{ fm}^{-1}$. In other words, the IK functions could not explain simultaneously the mass splitting of positive parity baryons and the proton radius. Usage of such nonrelativistic wave function to describe the Λ^* baryon, moreover an attempt to generalize and apply it to the relativistic multiquark system like the hexaquark is, at least, questionable. Existing problems of the IK model were emphasized already in the original paper [26] and repeated in Ref. [32], nevertheless in both of them these wave functions were employed to carry out numerical computations.

Calculations demonstrate that the matrix element $|\mathcal{M}(\Lambda^*\Lambda^* \rightarrow S)|^2$ is highly sensitive to the choice of the parameters α_B and α_S , and is suppressed if the hexaquark S has a small radius $r_S/r_N \ll 1$, where r_N is the nucleon radius. The radius of the hexaquark was estimated in Ref. [29] as $0.1 \leq r_S \leq 0.3 \text{ fm}$ which implies $r_S/r_N = 0.11 - 0.34$. Then, for example, for $r_S \approx 0.13 \text{ fm}$ and, as a result, for the ratio $r_N/r_S \approx 6.6$ the BBG wave function with the core radius $c \simeq 0.4 \text{ fm}$, and $\alpha_B = 0.406 \text{ fm}^{-1}$ satisfies the constraint $|\mathcal{M}(\Lambda^*\Lambda^* \rightarrow S)|^2 \lesssim 10^{-25}$ necessary to evade the SK bound (see, Fig. 1 in Ref. [26]). At the same parameters of the hexaquark, but for $\alpha_B = 0.221 \text{ fm}^{-1}$ the core radius in the BBG model should be larger to achieve required suppression of the matrix element $|\mathcal{M}(\Lambda^*\Lambda^* \rightarrow S)|^2$.

To update predictions for $|\mathcal{M}(\Lambda^*\Lambda^* \rightarrow S)|^2$, the authors in Ref. [32] used the new two-nucleon wave functions. The latter were extracted by utilizing an information on the two-nucleon point density inside nuclei $\rho_{NN}(r)$ [59]. The original calculations of $\rho_{NN}(r)$ were performed by employing the nonrelativistic Hamiltonian, where the phenomenological NN potential includes electromagnetic and one-pion-exchange terms, and

also contains phenomenological contributions to reproduce nucleon-nucleon elastic scattering. The new wave functions, of course, present a more detailed picture of a nucleus, but are nonrelativistic quantities and do not take into account dynamical effects of quark-gluon interactions in nuclei. The updated results for $|\mathcal{M}(\Lambda^*\Lambda^* \rightarrow S)|^2$ are presented in Fig. 5 of Ref. [32] as a function of the ratio r_S/r_N . Unfortunately, the authors did not show the region $r_S/r_N = 0.11 - 0.2$, in which the restriction on the matrix element $|\mathcal{M}(\Lambda^*\Lambda^* \rightarrow S)|^2 \lesssim 10^{-25}$ may be satisfied.

As it has been just emphasized above, $|\mathcal{M}(\Lambda^*\Lambda^* \rightarrow S)|^2$ critically depends on a behavior of the relevant wave functions at small inter-nucleon distances $r \lesssim 1 \text{ fm}$. At these distances nucleons are not the suitable degrees of freedom, and quark-gluon content of the nucleons becomes essential to describe correctly processes inside nuclei. Neither the Isgur-Karl type wave functions of the Λ^* baryon and S-hexaquark, nor the two-nucleon wave functions discussed till now contain detailed information on relativistic and nonperturbative interactions of quarks and gluons at distances $r \lesssim 1 \text{ fm}$ and high densities. Therefore, the predictions for $|\mathcal{M}(\Lambda^*\Lambda^* \rightarrow S)|^2$ made in Refs. [26, 32] cannot be considered as credible ones and used to confirm or exclude existence of the hexaquark S. Only after thorough exploration of aforementioned problems, we can get serious estimate for $|\mathcal{M}(\Lambda^*\Lambda^* \rightarrow S)|^2$ and see whether the requirement $r_S \ll r_N$ —necessary for stability of $^{16}\text{O}_8$ in the present picture—survives or not. Then we will be able to answer the question moved to the title of the present article as well.

Appendix: Details of calculations and the correlation function $\Pi^{\text{OPE}}(M^2, s_0)$

In this appendix we present some details of calculations, which are necessary to derive the sum rules for the mass and coupling of the hexaquark. It is evident that a main problem is calculation of the QCD side of the sum rules $\Pi^{\text{OPE}}(p)$. Using explicit expressions for the light quarks propagators and inserting relevant ones into Eq. (6), we get the Fourier integrals of following types:

$$T[l, m] = \int d^4x e^{ipx} \frac{[1, x_\mu, x_\mu x_\nu, \dots](x^2)^l \left[\ln \left(\frac{-x^2 \Lambda^2}{4} \right) \right]^m}{(x^2)^n}. \quad (\text{A.1})$$

For simplicity, let us consider the case $l = 0$ and $m = 0$. After a Wick rotation in the Euclidean space for $T[0, 0] \equiv T$, one finds

$$T = -i(-1)^n \int d^4x_E \frac{e^{-ip_E x_E}}{(x_E^2)^n}. \quad (\text{A.2})$$

By applying the Schwinger parametrization

$$\frac{1}{A^n} = \frac{1}{\Gamma(n)} \int_0^\infty dt t^{n-1} e^{-tA}, \quad A > 0, \quad (\text{A.3})$$

it is not difficult to recast T into the form

$$T = -i(-1)^n \frac{1}{\Gamma(n)} \int_0^\infty dt t^{n-1} \int d^4 x_E e^{-ip_E x_E} e^{-tx_E^2}.$$

By performing the resultant Gaussian integral over four- x , we obtain

$$T = -i(-1)^n \frac{\pi^2}{\Gamma(n)} \int_0^\infty dt t^{n-3} e^{-p_E^2/4t}. \quad (\text{A.4})$$

The next step is to apply the Borel transformation with respect to p_E^2 to suppress contributions of the higher resonances and continuum states. To this end, we utilize the formula

$$\mathcal{B}_M e^{-p_E^2/4t} = \delta\left(\frac{1}{M^2} - \frac{1}{4t}\right), \quad (\text{A.5})$$

which leads to

$$\mathcal{B}_M T = -i(-1)^n \frac{\pi^2}{\Gamma(n)} \int_0^\infty dt t^{n-3} \delta\left(\frac{1}{M^2} - \frac{1}{4t}\right). \quad (\text{A.6})$$

By carrying out the integration over t , we immediately get

$$\mathcal{B}_M T = -i(-1)^n \frac{\pi^2}{\Gamma(n)} \left(\frac{M^2}{4}\right)^{n-3}. \quad (\text{A.7})$$

Afterwards we apply the continuum subtraction procedure using the replacement

$$(M^2)^N \rightarrow \frac{1}{\Gamma(N)} \int_0^{s_0} ds e^{-s/M^2} s^{N-1}, \quad N > 0, \quad (\text{A.8})$$

where s_0 is the continuum threshold parameter. Then for the Borel transformed and subtracted integral T , we get

$$\mathcal{B}_M T = -i(-1)^n \frac{\pi^2 4^{3-n}}{\Gamma(n)\Gamma(n-3)} \int_0^{s_0} ds e^{-s/M^2} s^{N-1}, \quad (\text{A.9})$$

Calculations of the other terms $T[l, m]$ in QCD side of the sum rule can be performed in a similar manner. In the general case of $T[l, m]$, for continuum subtraction one should use more complicated formulas, full list of which can be found in Ref. [60]

As a result, for $\Pi^{\text{OPE}}(M^2, s_0)$ we get

$$\Pi^{\text{OPE}}(M^2, s_0) = \frac{1}{147 \cdot 4^{12} 5^2 \pi^{10}} \int_0^{s_0} ds s^7 e^{-s/M^2} + \sum_{i=3}^{30} \Pi^{(i)}(M^2, s_0), \quad (\text{A.10})$$

where the first term is the perturbative contribution.

For the nonperturbative (3) – (10) dimensional terms we get:

$$\begin{aligned}
\Pi^{(3)} &= \frac{7m_s \langle \bar{s}s \rangle}{3^3 4^9 5^2 \pi^8} \int_0^{s_0} ds s^5 e^{-s/M^2}, \\
\Pi^{(4)} &= \frac{\langle g_s^2 G^2 \rangle}{2 \cdot 4^{11} 5^2 \pi^{10}} \int_0^{s_0} ds s^5 e^{-s/M^2}, \\
\Pi^{(5)} &= \frac{m_s}{2 \cdot 3^3 4^{11} 5^2 \pi^8} [2 \langle \bar{q}q_s \sigma Gq \rangle \\
&\times \int_0^{s_0} ds s^4 [12783 - 2340 \ln(s/\Lambda^2)] e^{-s/M^2} + \langle \bar{s}g_s \sigma Gs \rangle \int_0^{s_0} ds s^4 [346 + 120 \ln(s/\Lambda^2)] e^{-s/M^2}], \\
\Pi^{(6)} &= \frac{1}{2 \cdot 3^3 4^7 5 \pi^6} [2 \langle \bar{q}q \rangle^2 + \langle \bar{s}s \rangle^2 + 12 (2 \langle \bar{s}s \rangle \langle \bar{q}q \rangle + \langle \bar{q}q \rangle^2)] \int_0^{s_0} ds s^4 e^{-s/M^2}, \\
\Pi^{(7)} &= \frac{m_s \langle g_s^2 G^2 \rangle}{3^3 4^{11} \pi^8} \left[5 \langle \bar{q}q \rangle \int_0^{s_0} ds s^3 [-311 + 60 \ln(s/\Lambda^2)] e^{-s/M^2} + \langle \bar{s}s \rangle \int_0^{s_0} ds s^3 [865 - 36 \ln(s/\Lambda^2)] e^{-s/M^2} \right], \\
\Pi^{(8)} &= -\frac{1}{2 \cdot 3^4 4^{15} \pi^{10}} [1429 \langle g_s^2 G^2 \rangle^2 + 21576 \pi^4 m_0^2 (181 \langle \bar{q}q \rangle^2 + 5 \langle \bar{s}s \rangle^2 + 342 \langle \bar{q}q \rangle \langle \bar{s}s \rangle)] \int_0^{s_0} ds s^3 e^{-s/M^2}, \\
\Pi^{(9)} &= \frac{m_s}{2 \cdot 3^4 4^{12} \pi^8} \langle g_s^2 G^2 \rangle m_0^2 \int_0^{s_0} ds s^2 [-77464 \langle \bar{q}q \rangle + 13922 \langle \bar{s}s \rangle - (27320 \langle \bar{q}q \rangle - 748 \langle \bar{s}s \rangle) \ln(s/\Lambda^2)] e^{-s/M^2} \\
&- \frac{m_s}{2 \cdot 3^3 4^4 \pi^4} [-12 \langle \bar{s}s \rangle^2 \langle \bar{q}q \rangle + 14 \langle \bar{s}s \rangle \langle \bar{q}q \rangle^2 + 24 \langle \bar{q}q \rangle^3] \int_0^{s_0} ds s^2 e^{-s/M^2}, \\
\Pi^{(10)} &= \frac{\langle g_s^2 G^2 \rangle}{3^3 4^9 \pi^6} [386 \langle \bar{q}q \rangle^2 + 34 \langle \bar{s}s \rangle^2 + 636 \langle \bar{s}s \rangle \langle \bar{q}q \rangle] \int_0^{s_0} ds s^2 e^{-s/M^2} - \frac{m_0^4}{3 \cdot 4^9 \pi^6} [27 \langle \bar{q}q \rangle^2 + \langle \bar{s}s \rangle^2] \int_0^{s_0} ds s^2 e^{-s/M^2}.
\end{aligned} \tag{A.11}$$

In $\Pi^{(i)}(M^2, s_0)$ we have assumed $\langle \bar{u}u \rangle = \langle \bar{d}d \rangle$ and denoted both of them as $\langle \bar{q}q \rangle$. We do not provide explicit expressions of the terms (i) > (10).

-
- | | |
|--|---|
| <p>[1] R. L. Jaffe, Phys. Rev. D 15, 267 (1977).
 [2] H. X. Chen, W. Chen, X. Liu and S. L. Zhu, Phys. Rept. 639, 1 (2016).
 [3] H. X. Chen, W. Chen, X. Liu, Y. R. Liu and S. L. Zhu, Rept. Prog. Phys. 80, 076201 (2017).
 [4] A. Esposito, A. Pilloni and A. D. Polosa, Phys. Rept. 668, 1 (2017).
 [5] A. Ali, J. S. Lange and S. Stone, Prog. Part. Nucl. Phys. 97, 123 (2017).
 [6] S. L. Olsen, T. Skwarnicki and D. Zieminska, Rev. Mod. Phys. 90, 015003 (2018).
 [7] R. L. Jaffe, Phys. Rev. Lett. 38, 195 (1977) Erratum: [Phys. Rev. Lett. 38, 617 (1977)].
 [8] K. F. Liu and C. W. Wong, Phys. Lett. 113B, 1 (1982).
 [9] S. A. Yost and C. R. Nappi, Phys. Rev. D 32, 816 (1985).
 [10] M. Oka, K. Shimizu and K. Yazaki, Phys. Lett. 130B, 365 (1983).
 [11] A. P. Balachandran, F. Lizzi, V. G. J. Rodgers and A. Stern, Nucl. Phys. B 256, 525 (1985).
 [12] M. Oka, K. Shimizu and K. Yazaki, Nucl. Phys. A 464, 700 (1987).
 [13] U. Straub, Z. Y. Zhang, K. Brauer, A. Faessler and S. B. Khadkikar, Phys. Lett. B 200, 241 (1988).
 [14] Y. Koike, K. Shimizu and K. Yazaki, Nucl. Phys. A 513, 653 (1990).</p> | <p>[15] S. A. Larin, V. A. Matveev, A. A. Ovchinnikov and A. A. Pivovarov, Sov. J. Nucl. Phys. 44, 690 (1986) [Yad. Fiz. 44, 1066 (1986)] [hep-ph/0405035].
 [16] N. Kodama, M. Oka and T. Hatsuda, Nucl. Phys. A 580, 445 (1994).
 [17] Y. Iwasaki, T. Yoshie and Y. Tsuboi, Phys. Rev. Lett. 60, 1371 (1988).
 [18] S. R. Beane <i>et al.</i> [NPLQCD Collaboration], Phys. Rev. D 87, 034506 (2013).
 [19] A. Francis, J. R. Green, P. M. Junnarkar, C. Miao, T. D. Rae and H. Wittig, Phys. Rev. D 99, 074505 (2019).
 [20] H. Suganuma and K. Matsumoto, EPJ Web Conf. 137, 13018 (2017).
 [21] A. Alavi-Harati <i>et al.</i> [KTeV Collaboration], Phys. Rev. Lett. 84, 2593 (2000).
 [22] T. F. Carames and A. Valcarce, Int. J. Mod. Phys. E 22, 1330004 (2013).
 [23] B. H. Kim <i>et al.</i> [Belle Collaboration], Phys. Rev. Lett. 110, 222002 (2013).
 [24] J. P. Lees <i>et al.</i> [BaBar Collaboration], Phys. Rev. Lett. 122, 072002 (2019).
 [25] G. R. Farrar and G. Zaharijas, Phys. Lett. B 559, 223 (2003). Erratum:[Phys. Lett. B 575, 358 (2003).]
 [26] G. R. Farrar and G. Zaharijas, Phys. Rev. D 70, 014008</p> |
|--|---|

- (2004).
- [27] G. R. Farrar, arXiv:1708.08951 [hep-ph].
- [28] G. R. Farrar, PoS ICRC **2017**, 929 (2018).
- [29] G. R. Farrar, arXiv:1805.03723 [hep-ph].
- [30] E. W. Kolb and M. S. Turner, Phys. Rev. D **99**, 063519 (2019).
- [31] S. D. McDermott, S. Reddy and S. Sen, Phys. Rev. D **99**, 035013 (2019).
- [32] C. Gross, A. Polosa, A. Strumia, A. Urbano and W. Xue, Phys. Rev. D **98**, 063005 (2018).
- [33] L. Maiani, F. Piccini, A. D. Polosa, and V. Riquer, Phys. Rev. Lett. **93**, 212002 (2004).
- [34] R. L. Jaffe, Phys. Rept. **409**, 1 (2005).
- [35] M. Tanabashi *et al.* (Particle Data Group), Phys. Rev. D **98**, 030001 (2018).
- [36] E. Aprile *et al.* [XENON Collaboration], Phys. Rev. Lett. **123**, 251801 (2019).
- [37] E. Aprile *et al.* [XENON Collaboration], Phys. Rev. Lett. **123**, 241803 (2019).
- [38] R. Agneze *et al.* [SuperCDMS Collaboration], Phys. Rev. D **99**, 062001 (2019).
- [39] M. Shafi Mahdawi, and G. R. Farrar, JCAP **12**, 004 (2017).
- [40] D. McCammon *et al.*, Astrophys. J. **576**, 188 (2002).
- [41] G. Bertone, D. Hooper, and J. Silk, Phys. Rept. **405**, 279 (2005).
- [42] H. Kim, K. S. Kim, M. K. Cheoun and M. Oka, Phys. Rev. D **97**, 094005 (2018).
- [43] S. S. Agaev, K. Azizi and H. Sundu, Phys. Lett. B **781**, 279 (2018).
- [44] S. S. Agaev, K. Azizi and H. Sundu, Phys. Lett. B **784**, 266 (2018).
- [45] W. Ochs, J. Phys. G. **40**,043001 (2013).
- [46] L. S. Kisslinger, and W. H. Ma, Phys. Lett. B **485**, 367 (2000).
- [47] G. Erkol, R. G. E. Timmermans, and Th. A. Rijken, Phys. Rev. C **72**, 035209 (2005).
- [48] D. Hooper, and C. D. McDermott, Phys. Rev. D **97**, 115006 (2018).
- [49] O. Adriani *et al.* [PAMELA Collaboration], Riv. Nuovo Cim. **40**, 1 (2017).
- [50] M. Aguilar *et al.* [AMS Collaboration], Phys. Rev. Lett. **110**, 141102 (2013).
- [51] O. Adriani *et al.* [PAMELA Collaboration], Phys. Rev. Lett. **105**, 121101 (2010).
- [52] M. Aguilar *et al.* [AMS Collaboration], Phys. Rev. Lett. **116**, 091103 (2016).
- [53] F. A. Aharonian, A. M. Atoyan, and H. J. Volk, Astron. Astrophys. **294**, L41 (1995).
- [54] P. Blasi, Phys. Rev. Lett. **103**, 051104 (2009).
- [55] M. Ahlers, P. Mertsch, and S. Sarkar, Phys. Rev. D **80**, 123017 (2009).
- [56] D. Clowe, A. Gonzalez, and M. Markevitch, Astrophys. J. **604**, 596 (2004).
- [57] D. Wadekar, and G. R. Farrar, arXiv:1903.12190 [hep-ph].
- [58] N. Isgur and G. Karl, Phys. Rev. D **19**, 2653 (1979). Erratum: [Phys. Rev. D **23**, 817 (1981).]
- [59] D. Lonardonì, A. Lovato, S. C. Pieper and R. B. Wiringa, Phys. Rev. C **96**, 024326 (2017).
- [60] S. S. Agaev, K. Azizi and H. Sundu, Phys. Rev. D **93**, 114036 (2016).