Nuclear-Bound Quarkonium

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We show that the QCD van der Waals interaction due to multiple-gluon exchange provides a new kind of attractive nuclear force capable of binding heavy quarkonia to nuclei. The parameters of the potential are estimated by identifying multigluon exchange with the Pomeron contributions to elastic meson-nucleon scattering. The gluonic potential is then used to study the properties of $c\bar{c}$ nuclear-bound states. In particular, we predict bound states of the η_c with ^{3}He and heavier nuclei. Production modes and rates are also discussed.

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One of the most interesting anomalies in hadron physics is the remarkable behavior of the spin-spin correlation A_{NN} for $pp \to pp$ elastic scattering at $\theta_{\mathrm{c.m.}} = 90°$: As $\sqrt{s}$ crosses 5 GeV the ratio of cross sections for protons scattering with their incident spins parallel and normal to the scattering plane and for protons scattering with their spins antiparallel changes rapidly from approximately 2:1 to 4:1.[1] As shown in Ref. 2, this behavior can be understood as the consequence of a strong threshold enhancement at the open-charm threshold for $pp \to \Lambda_c D p$ at $\sqrt{s} = 5.08$ GeV. The dominant enhancement in the $pp \to pp$ amplitude occurs in the partial wave $J = L = S = 1$, which matches the quantum numbers of the $J = 1$ S-wave eight-quark system $qqqqqq(c\bar{c})_{S=1}$ at threshold. Strong final-state interactions are expected at the threshold for new-flavor production, since at threshold, all the quarks in the final state have nearly zero relative velocity.

In this paper we discuss the possibility of production of *hidden* charm below threshold in hadronic and nuclear collisions. Consider the reaction $pd \to (c\bar{c})^3$He where the charmonium state is produced nearly at rest. At the threshold for charm production, the incident nuclei will be nearly stopped (in the center-of-mass frame) and will fuse into a compound nucleus (the ^{3}He) because of the strong attractive nuclear force. The charmonium state will be attracted to the nucleus by the QCD gluonic van der Waals force. One thus expects strong final-state interactions near threshold. In fact, we shall argue that the $c\bar{c}$ system will bind to the ^{3}He nucleus. It is thus likely that a new type of exotic nuclear bound state will be formed: charmonium bound to nuclear matter. Such a state should be observable at a distinct pd energy, spread by the width of the charmonium state, and it will decay to unique signatures such as $pd \to {}^3\mathrm{He}\gamma\gamma$. The binding energy in the nucleus gives a measure of the charmonium's interactions with ordinary hadrons and nuclei; its decays will measure hadron-nucleus interactions and test color transparency starting from a unique initial-state condition.

In quantum chromodynamics, a heavy-quarkonium $Q\bar{Q}$ state such as the η_c interacts with a nucleon or nucleus through multiple gluon exchange. This is the QCD analog of the attractive QED van der Waals potential. Unlike QED, the potential cannot have an inverse power law at large distances because of the absence of zero-mass gluonium states. Since the $(Q\bar{Q})$ and nucleons have no quarks in common, the quark interchange (or equivalently the effective meson-exchange) potential should be negligible. Since there is no Pauli blocking, the effective quarkonium-nuclear interaction will not have a short-range repulsion.

The QCD van der Waals interaction is the simplest example of a nuclear force in QCD.[3] In this paper we shall show that this potential is sufficiently strong to bind quarkonium states such as the η_c to nuclear matter. The signal for such states will be narrow peaks in energy in the production cross section. On general grounds one expects that the effective nonrelativistic potential between heavy quarkonium and nucleons can be parametrized by a Yukawa form $V_{(Q\bar{Q})A} = -\alpha e^{-\mu r}/r$. Since the gluons have spin 1, the interaction is vectorlike. This implies a rich spectrum of quarkonium-nucleus bound states with spin-orbit and spin-spin hyperfine splitting.

Thus far lattice gauge theory and other nonperturbative methods have not determined the range or magnitude of the gluonic potential between hadrons. A phenomenological estimate of the effective Yukawa potential applicable to $Q\bar{Q}Q\bar{Q}$ bound states and resonances has been given by Liu.[4] Using experimental evidence for resonant structure in both ϕ-ϕ and ψ-ψ systems,[5] he finds an effective Yukawa potential for ϕ-ψ interactions with strength $\alpha \sim 0.45$ and inverse range $\mu = 1$ GeV. One would expect the η_c-N interaction to have a somewhat larger magnitude and longer range.

We can obtain another estimate of the effective QCD van der Waals potential using high-energy information. Here we have in mind a QED analog, where one may use high-energy lepton-lepton scattering to normalize the Coulomb potential at low energies. In our case, we relate vector-exchange contributions to the high-energy forward hadron-nucleon scattering amplitude to the vector part of the multigluon-exchange potential at low energies. This should be a good approximation to the extent that the QCD van der Waals potential behaves like a local spin-1 potential, analogous to vector-boson exchange. In effect we identify Pomeron exchange with the eikonalization of the two-gluon-exchange potential.

We shall use the phenomenological model of high-energy Pomeron interactions developed by Donnachie and Landshoff,[6] where, as in the Chou-Yang model, one has an s-independent parametrization of the meson-nucleon and meson-nucleus cross sections at small t:

$$d\sigma/dt(MA \to MA) = [2\beta F_M(t)]^2[3A\beta F_A(t)]^2/4\pi .$$

Here $\beta = 1.85$ GeV^{-1} is the flavor-independent Pomeron-quark coupling constant, and A is the nucleon number of the nucleus. To first approximation the form factors can be identified with the helicity-zero meson and nuclear electromagnetic form factors.

Ignoring corrections due to eikonalization, the cross section at $s \gg |t|$ can be identified with that due to the vector Yukawa potential $d\sigma/dt(MA \to MA) = 4\pi\alpha^2/(-t+\mu^2)^2$. We then can calculate the effective coupling α and the range μ from $(d\sigma/dt)^{1/2}$ and its slope at $t=0$. For π-N or K-N scattering, one obtains $\alpha \simeq 0.5$ and $\mu \simeq 0.5$ GeV.

In the case of the η_c-N interaction, the radius of the meson is roughly half that of ordinary mesons.[7] In the fully screened model of Gunion and Soper,[8] this reduces the total cross section by a factor of 4. In the Donnachie-Landshoff model,[6] screening is not so effective because of the approximately local coupling of the Pomeron to hadrons. If we assume partial screening corresponding to a reduction of the elastic cross section by a factor of 2 relative to K-N scattering,[9] we obtain $\alpha = 0.4$ and $\mu = 0.6$ GeV for the effective η_c-N potential. Nonscreening corresponds to $\alpha = 0.6$ and $\mu = 0.6$ GeV. In either case these estimates are roughly consistent with Liu's values.

In the case of η_c-nucleus scattering, the slope is dominated by the nuclear size since the $(c\bar{c})$ radius is comparatively small; thus $\mu^{-2} = |dF_A(t)/dt|_{t=0} = \langle R_A^2 \rangle/6$ and $\alpha = 3A\beta^2\mu^2/2\pi$, for the unscreened case. For η_c-^{3}He scattering, one finds $\alpha \simeq 0.3$ and the inverse range is $\mu \simeq 250$ MeV reflecting the smearing of the local interaction over the nuclear volume. For partially screened charmonium interactions, α is reduced by $1/\sqrt{2}$. The QCD van der Waals potential is effectively the only QCD interaction between the charmonium bound state and nuclear matter. In the threshold regime the η_c is nonrelativistic, and an effective-potential Schrödinger equation of motion is applicable. To first approximation we will treat the η_c as a stable particle. The effective potential is then real since higher-energy intermediate states from charmonium or nuclear excitations should not be important.

We compute the binding energy using the variational wave function $\psi(r) = (\gamma^3/\pi)^{1/2}\exp(-\gamma r)$. Results for binding energies are presented in Table I for values of the coupling constant corresponding to the unscreened and partially screened cases. The condition for binding by the Yukawa potential with the above wave function is $\alpha m_{\text{red}} > \mu$. This condition is not met for charmonium-proton or charmonium-deuterium systems. However, the binding of the η_c to a heavy nucleus increases rapidly with A, since the potential strength is linear in A, and the kinetic energy $\langle \mathbf{p}^2/2m_{\text{red}} \rangle$ decreases faster than the square of the nuclear size. If the width of the $c\bar{c}$ is much smaller than its binding energy, the charmonium state lives sufficiently long that it can be considered stable for the purposes of calculating its binding to the nucleus. For η_c ^{3}He and $\alpha = 0.327$ the computed binding energy is 19 MeV, and for η_c ^{4}He the binding energy is over 140 MeV. The predicted binding energies are large even though the QCD van der Waals potential is relatively weak compared to the one-pion-exchange Yukawa potential; this is due to the absence of Pauli blocking or a repulsive short-range potential for heavy quarks in the nucleus. Table I gives a list of computed binding energies for $c\bar{c}$-nuclear systems. A two-parameter variational wave function of the form $(e^{-\alpha_1 r} - e^{-\alpha_2 r})/r$ gives essentially the same results. Our results also have implications for the binding of hidden strange hadrons to nuclei.[11] However, the strong mixing of the η with nonstrange quarks makes the interpretation of such states

TABLE I. Binding energies $\epsilon = |\langle H \rangle|$ of the η_c to various nuclei, for unscreened and partially screened potentials; all masses are given in GeV. The data for $\langle R_A^2 \rangle^{1/2}$ (in GeV^{-1}) are from Ref. 10.

				Unscreened			Partially screened		
A	$\langle R_A^2 \rangle^{1/2}$	μ	m_{red}	α	γ	$\langle H \rangle$	α	γ	$\langle H \rangle$
1	3.9	0.60	0.715	0.59		>0	0.42		>0
2	10.7	0.23	1.15	0.172		>0	0.122		>0
3	9.5	0.26	1.45	0.327	0.40	-0.019	0.231	0.24	-0.003
4	8.2	0.30	1.66	0.585	0.92	-0.143	0.414	0.62	-0.051
6	11.2	0.22	1.95	0.470	0.89	-0.128	0.332	0.61	-0.050
9	11.2	0.22	2.20	0.705	1.53	-0.407	0.499	1.07	-0.179

more complicated.

It is clear that the production cross section for charm production near threshold in nuclei will be very small. We estimate rates below. However, the signals for bound $c\bar{c}$ to nuclei are very distinct. The most practical measurement could be the inclusive process $pd \to {}^3\mathrm{He}X$, where the missing mass M_X is constrained close to the charmonium mass. Since the decay of the bound $c\bar{c}$ is isotropic in the center-of-mass frame, but backgrounds are peaked forward, the most favorable signal to noise is at backward ^{3}He c.m. angles. If the η_c is bound to the ^{3}He, a peak will be found at a distinct value of incident pd energy: $\sqrt{s} = M_{\eta_c} + M_{{}^3\mathrm{He}} - \epsilon$, spread by the intrinsic width of the η_c. Here ϵ is the η_c-nucleus binding energy predicted from the Schrödinger equation.

The momentum distribution of the outgoing nucleus in the center-of-mass frame is given by $dN/d^3p = |\psi(\mathbf{p})|^2$. Thus the momentum distribution gives a direct measure of the $c\bar{c}$-nuclear wave function. The width of the momentum distribution is given by the wave-function parameter γ, which is tabulated in Table I. The kinematics for several different reactions are given in Table II. From the uncertainty principle we expect that the final-state momentum $\mathbf{p}$ is related inversely to the uncertainty in the $c\bar{c}$ position when it decays. By measuring the binding energy and recoil momentum distribution in $\mathbf{p}$, one determines the Schrödinger wave function, which then can be easily inverted to give the quarkonium-nuclear potential.

Energy conservation in the center-of-mass frame implies $E_{\mathrm{c.m.}} = E_X + E_A \simeq M_X + M_A + \mathbf{p}^2/2M_r'$. Here $M_r' = (1/M_X + 1/M_A)^{-1}$ is the reduced mass of the final-state system. The missing invariant mass is always less than the mass of the free η_c: $M_X = M_{\eta_c} - \epsilon - \mathbf{p}^2/2M_r'$; thus the invariant mass varies with the recoil momentum. The mass deficit can be understood as the result of the fact that the η_c decays off its energy shell when bound to the nucleus.

More information is obtained by studying completely specified final states—exclusive channels such as $pd \to \gamma\gamma\,{}^3\mathrm{He}$. Observation of the two-photon decay of the η_c would be a decisive signal for nuclear-bound quarkonia. The position of the bound $c\bar{c}$ at the instant of its decay is distributed in the nuclear volume according to the eigen wave function $\psi(\mathbf{r})$. Thus the hadronic decays of the $c\bar{c}$ system allow the study of the propagation of hadrons through the nucleus starting from a wave packet centered on the nucleus, a novel initial condition. In each case, the initial-state condition for the decay is specified by the Schrödinger wave function with specific orbital and spin quantum numbers. Consider, then, the decay $\eta_c \to p\bar{p}$. As the nucleons transit the nuclear medium, their outgoing wave will be modified by nuclear final-state interactions. The differential between the energy and momentum spectrum of the proton and antiproton should be a sensitive measure of the hadronic amplitudes. More interesting is the fact that the nucleons are initially formed from the $c\bar{c} \to gg$ decay amplitude. The size of the production region is of the order of the charm Compton length $l \sim 1/m_c$. The proton and antiproton thus interact in the nucleus as a small color-singlet state before they are asymptotic hadron states. The distortion of the outgoing hadron momenta thus tests formation-zone physics[9] and color transparency.[12]

The interactions of the J/ψ and other excited states of charmonium in nuclear matter are more complicated than the η_c interaction because of the possibility of spin-exchange interactions which allow the $c\bar{c}$ system to couple to the η_c. This effect adds inelasticity to the effective $c\bar{c}$ nuclear potential. In effect the bound J/ψ-^{3}He can decay to $\eta_c dp$ and its width will change from tens of keV to MeV. However, if the J/ψ-nucleus binding is sufficiently strong, then the η_c plus nuclear continuum states may not be allowed kinematically, and the bound J/ψ could then retain its narrow width, $\simeq 70$ keV. As seen in Table I this appears to be the case for the J/ψ-^{4}He system in the unscreened case. An important signature for the bound vector charmonium state will be the exclusive l^+l^- plus nucleus final state.

The narrowness of the charmonium states implies that the charmonium-nucleus bound state is formed at a sharp distinct c.m. energy, spread by the total width Γ and the much smaller probability that it will decay back to the initial state. By duality the product of the cross-section peak times its width should be roughly a constant. Thus the narrowness of the resonant energy leads to a large multiple of the peak cross section, favoring experiments with good incident-energy resolution.

The formation cross section is thus characterized by a series of narrow spikes corresponding to the binding of the various $c\bar{c}$ states. In principle, there could be higher-orbital or higher-angular-momentum bound-state excitations of the quarkonium-nuclear system. In the case of J/ψ bound to spin-half nuclei, we predict a hyperfine separation of the $L=0$ ground state corresponding to states of total spin $J=\frac{3}{2}$ and $J=\frac{1}{2}$. This separation will measure the gluonic magnetic moment of the nucleus and that of the J/ψ. Measurements of the binding energies could, in principle, be done with excellent precision, thus determining fundamental hadronic measures with high accuracy.

The production cross section for creating the quark-

TABLE II. Kinematics for the production of η_c-nucleus bound states. All quantities are given in GeV.

Process	ϵ	$p_{\mathrm{c.m.}}$	p_1^{lab}
$\gamma\,{}^3\mathrm{He} \to ({}^3\mathrm{He}\,\eta_c)$	0.020	2.20	4.52
$pd \to ({}^3\mathrm{He}\,\eta_c)$	0.020	2.48	7.64
$\bar{p}\,{}^4\mathrm{He} \to ({}^3\mathrm{H}\,\eta_c)$	0.020	1.48	2.30
$\gamma\,{}^4\mathrm{He} \to ({}^4\mathrm{He}\,\eta_c)$	0.120	2.24	3.96
$n\,{}^3\mathrm{He} \to ({}^4\mathrm{He}\,\eta_c)$	0.120	2.60	6.09
$dd \to ({}^4\mathrm{He}\,\eta_c)$	0.120	2.71	9.51

onium-nucleus bound state is suppressed by a dynamical "stopping" factor representing the probability that the nucleons and nuclei in the final state convert their kinetic energy to the heavy-quark pair and are all brought to approximately zero relative velocity. For example, in the reaction $pd \rightarrow (c\bar{c})^3\mathrm{He}$ the initial proton and deuteron must each change momentum from $p_{\mathrm{c.m.}}$ to zero momentum in the center-of-mass frame. The probability for a nucleon or nucleus to change momentum and stay intact is given by the square of its form factor $F_A^2(q_A^2)$, where $q_A^2 = [(M_A^2 + p_{\mathrm{c.m.}}^2)^{1/2} - M_A]^2 - p_{\mathrm{c.m.}}^2$. We can use as a reference cross section the $pp \rightarrow c\bar{c}pp$ cross section above threshold, which was estimated in Ref. 2 to be of order $\sim 1\ \mu\mathrm{b}$. Then

$$\sigma(A_1A_2 \rightarrow c\bar{c}A_{12}) = \sigma(pp \rightarrow c\bar{c}pp)\frac{F_{A_1}^2(q_{A_1}^2)F_{A_2}^2(q_{A_2}^2)}{F_N^4(q_N^2)}\,. \tag{1}$$

For the $pd \rightarrow (c\bar{c})^3\mathrm{He}$ channel, we thus obtain a suppression factor relative to the pp channel of $F_d^2(4.6\ \mathrm{GeV}^2)F_p^2(3.2\ \mathrm{GeV}^2)/F_N^4(2.8\ \mathrm{GeV}^2) \sim 10^{-5}$ giving a cross section which may be as large as $10^{-35}\ \mathrm{cm}^2$. Considering the uniqueness of the signal and the extra enhancement at the resonance energy, this appears to be a viable experimental cross section.

In QCD, the nuclear forces are identified with the residual strong color interactions due to quark interchange and multiple-gluon exchange.[13] Because of the identity of the quark constituents of nucleons, a short-range repulsive component is also present (Pauli blocking). From this perspective, the study of heavy-quarkonium interactions in nuclear matter is particularly interesting: Because of the distinct flavors of the quarks involved in the quarkonium-nucleon interaction there is no quark exchange to first order in elastic processes, and thus no one-meson-exchange potential from which to build a standard nuclear potential. For the same reason, there is no Pauli blocking and consequently no short-range nuclear repulsion. The nuclear interaction in this case is purely gluonic and thus of a different nature from the usual nuclear forces.

We have discussed the signals for recognizing quarkonium bound in nuclei. The production of nuclear-bound quarkonium would be the first realization of hadronic nuclei with exotic components bound by a purely gluonic potential. It should be emphasized that the values assumed in this paper for the coupling and range of the charmonium-nuclear potential are only rough estimates. The measurement of η_c-nucleus binding energy would provide precise values for this interaction. Furthermore, the charmonium-nucleon interaction would provide the dynamical basis for understanding the spin-spin correlation anomaly in high-energy p-p elastic scattering.[2] In this case, the interaction is not strong enough to produce a bound state, but it can provide a strong enough enhancement at the heavy-quark threshold characteristic of an almost bound system.[14]

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[14]The signal for the production of almost bound nucleon (or nuclear) charmonium systems near threshold such as in $\gamma p \rightarrow (c\bar{c})p$ is the isotropic production of the recoil nucleon (or nucleus) at large invariant mass $M_X \cong M_{\eta_c}, M_{J/\psi}$.