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Nuclear-Bound Quarkonium*

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ABSTRACT

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 multi-gluon 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 He^3 and heavier nuclei. Production modes and rates are also discussed.

1. Introduction

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 $\theta_{cm} = 90°$: 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 to scattering with their spins anti-parallel 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 $\sqrt{s} = 5.08\ GeV$

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. The dominant enhancement in the $pp \to pp$ amplitude is expected 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, since the c and $\bar{c}$ have opposite parity. Even though the charm production rate is small, of order of $1\ \mu b$, it can have a large effect on the elastic $pp \to pp$ amplitude at $90°$ since the competing perturbative QCD hard-scattering amplitude at large momentum transfer is also very small at $\sqrt{s} = 5\ GeV$.

In this paper we discuss the possibility of production of **hidden** charm below threshold in hadronic and nuclear collisions. Consider the reaction $pd \to (c\bar{c})He^3$ 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 He^3) 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 He^3 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 $pd \to He^3\gamma\gamma$. 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.

2. The QCD van der Waals Interaction

In quantum chromodynamics, a heavy quarkonium $Q\overline{Q}$ state such as the η_c interacts with a nucleon or nucleus through multiple gluon exchange. This is the QCD analogue 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.[3] Since the $(Q\overline{Q})$ 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. In this paper we shall show that this potential is sufficiently strong to bind quarkonium states such as the η_c and η_b 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 non-relativistic potential between heavy quarkonium and nucleons can be parameterized by a Yukawa form

$$V_{(Q\overline{Q})A} = -\frac{\alpha e^{-\mu r}}{r}. \qquad (1)$$

Since the gluons have spin-one, the interaction is vector-like. 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 non-perturbative methods have not determined the range or magnitude of the gluonic potential between hadrons. However, we can obtain some constraint on the $J = 1$ flavor singlet interactions of hadrons by identifying the potential with the magnitude of the term linear in s in the meson-nucleon or meson-nucleus scattering amplitude. One can identify pomeron exchange with the eikonalization of the two-gluon exchange potential.[4]

To obtain a specific parameterization we shall make use of the phenomenological model of pomeron interactions developed by Donnachie and Landshoff.[5] These authors note that in order to account for the additive quark rule for total cross sections, the pomeron must have a somewhat local structure; its couplings are analogous to that of a heavy photon. The short-range character of the pomeron reflects the fact that the minimum gluonium mass which can be exchanged in the t-channel is of order several GeV. Interference terms between amplitudes involving different quarks can then be neglected.

The Donnachie-Landshoff formalism leads to an $s-$independent Chou-Yang parameterization of the meson-nucleon and meson-nucleus cross sections at small t:

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

Here $\beta = 1.85\ GeV^{-1}$ is the 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. We assume that β is independent of the meson type and nucleus. This is reasonable even for the η_C or J/ψ since the radius of the lowest charmonium state is $\simeq 0.5\ fm$, not very different than the pion radius.

Equation (2) gives a reasonable parameterization of the s-independent elastic hadron-hadron and hadron-nucleus scattering cross sections from very low to very high energies. Ignoring corrections due to eikonalization, we can identify the cross section at $s >> |t|$ with that due to the vector Yukawa potential

$$\frac{d\sigma}{dt}(MA \to MA) = \frac{4\pi\alpha^2}{(-t+\mu^2)^2}. \tag{3}$$

We calculate the effective coupling α and the range μ from $(d\sigma/dt)^{1/2}$ and its slope at $t = 0$. Assuming $0.5\ fm$ for the J/ψ radius, one obtains $\mu = 0.53\ GeV$, $\alpha = 0.46$, and an integrated elastic $J/\psi - N$ cross section $\sim 4\ mb.$[6] In the case of meson-nucleus scattering, the slope is dominated by the nuclear size since the $(c\bar{c})$ radius is comparatively small; thus $\mu^{-2} = |dF_A(t)/dt|_{t=0} =< R_A^2 > /6$ and $\alpha = 3A\beta^2\mu^2/2\pi$. For meson He^3 scattering, one finds $\alpha \simeq 0.3$ and $\mu \simeq 250 MeV$ reflecting the smearing of the local interaction over the nuclear volume.

In the case of η_c nucleus interactions, the QCD van der Waals potential is effectively the only QCD interaction. In the threshold regime the η_c is non-relativistic, 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 wavefunction $\psi(r) = \sqrt{(\gamma^3/\pi)}exp(-\gamma r)$. The condition for binding by the Yukawa potential with this wavefunction is $\alpha m_{red} > \mu$. 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 $< \vec{p}^2/2m_{red} >$ 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 $\eta_c He^3$ the computed binding energy is $\sim 20\ MeV$, and for $\eta_c He^4$ the binding energy is over $100\ MeV$. 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 the $c\bar{c}$ and $b\bar{b}$ nuclear systems. A two-parameter variational wavefunction 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 strange hadrons to nuclei.[7] However, the strong mixing of the η with non-strange quarks makes the interpretation of such states more complicated.

3. Searching for $c\bar{c}$ Nuclear-Bound States

It is clear that the production cross section for charm production near threshold in nuclei will be very small. We estimate rates in section 4. However the signals for bound $c\bar{c}$ to nuclei are very distinct. The most practical measurement could be the inclusive process $pd \rightarrow He^3 X$, where the missing mass M_X is constrained close to the charmonium mass. (See Fig. 1.) Since the decay of the bound $c\bar{c}$ is isotropic in the center-of-mass, but backgrounds are peaked forward, the most favorable signal-to-noise is at backward He^3 cm angles. If the η_c is bound to the He^3, a peak will be found at a distinct value of incident pd energy: $\sqrt{s} = M_{\eta_c} + M_{He^3} - \epsilon$, 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 $dN/d^3p = |\psi(\vec{p})|^2$. Thus the momentum distribution gives a direct measure of the $c\bar{c}$-nuclear wavefunction. The width of the momentum distribution is given by the wavefunction 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 $\vec{p}$ 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 $\vec{p}$, one determines the Schrödinger wavefunction, which then can be easily inverted to give the quarkonium-nuclear potential.

Energy conservation in the center of mass implies

$$E_{cm} = E_X + E_A \simeq M_X + M_A + \frac{\vec{p}^2}{2M_r'}. \tag{4}$$

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 :

$$M_X = M_{\eta_c} - \epsilon - \frac{\vec{p}^2}{2M_r'}; \tag{5}$$

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 $pd \to \gamma\gamma\ He^3$. 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-wavefunction $\psi(\vec{r})$. Thus the hadronic decays of the $c\bar{c}$ system allows 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 wavefunction 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 anti-proton 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 $\ell \sim 1/m_c$. The proton and anti-proton 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[6] and color transparency.[8]

4. Possibility of J/ψ-Nucleus Bound States

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, illustrated in Fig. 2, adds inelasticity to the effective $c\bar{c}$ nuclear potential. In effect the bound $J/\psi - He^3$ can decay to $\eta_c\ d\ p$ 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 $J/\psi - He^4$ system. An important signature for the bound vector charmonium state will be the exclusive $\ell^+\ell^-$ plus nucleus final state.

The narrowness of the charmonium states implies that the charmonium-nucleus bound state is formed at a sharp distinct cm 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 $J = 3/2$ and $J = 1/2$. 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.

5. Stopping Factor

The production cross section for creating the quarkonium-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 $pd \to (c\bar{c})He^3$ the initial proton and deuteron must each change momentum from $p_{\rm cm}$ to zero momentum in the center of mass. 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 $q_A^2 = [(M_A^2 + p_{\rm cm}^2)^{1/2} - M_A)^2 - p_{\rm cm}^2]$. We can use as a reference cross section the $pp \to c\bar{c}pp$ cross section above threshold, which was estimated in ref. 2 to be of order $\sim 1\ \mu b$. Then

$$\sigma(A_1A_2 \to c\bar{c}A_{12}) = \sigma(pp \to c\bar{c}pp)\frac{F_{A_1}^2(q_{A_1}^2)F_{A_2}^2(q_{A_2}^2)}{F_N^4(q_N^2)}. \tag{6}$$

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

6. Conclusions

In QCD, the nuclear forces are identified with the residual strong color interactions due to quark interchange and multiple-gluon exchange.[11] 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: due to 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. 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.[12]

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A	$\langle R_A^2\rangle^{1/2}$	μ	α	$(c\bar{c})$ m_{red}	γ	$\langle H\rangle$	$(b\bar{b})$ m_{red}	γ	$\langle H\rangle$
1	3.9	0.529	0.458	0.715		>0	0.85		>0
2	10.7	0.229	0.172	1.15		>0	1.563	0.18	–0.0012
3	9.5	0.26	0.327	1.45	0.40	–0.019	2.16	0.65	–0.050
	9.9	0.25	0.301		0.37	–0.015		0.60	–0.040
4	8.2	0.299	0.585	1.66	0.92	–0.143	2.66	1.52	–0.303
	8.4	0.292	0.557		0.87	–0.127		1.45	–0.271
	8.7	0.282	0.519		0.81	–0.107		1.35	–0.232
6	11.2	0.22	0.470	1.95	0.89	–0.128	3.50	1.63	–0.293
9	11.2	0.22	0.705	2.20	1.53	–0.407	4.42	3.11	–0.951
12	12.0	0.204	0.819	2.36	1.92	–0.637	5.09	4.16	–1.546
16	13.4	0.183	0.876	2.49	2.17	–0.805	5.74	5.0	–2.046

Table I

Binding energies $\epsilon = |<H>|$ of the η_c and η_b to various nuclei, in GeV. Here γ (in GeV) is the range parameter of the variational wavefunction, and μ (in GeV) and α are the parameters of the Yukawa potential. The data for $< R_A^2 >^{1/2}$ (in GeV^{-1}) are from ref. 9. We have assumed $M_{\eta_b} = 9.34\ GeV$.[10]

Process	M_1	M_2	M_A	ϵ	$\sqrt{s}$	p_{cm}	p_1^{lab}
$\gamma He^3 \to (He^3\eta_c)$	0	2.808	2.808	0.020	5.77	2.20	4.52
$pd \to (He^3\eta_c)$	0.938	1.876	2.808	0.020	5.77	2.48	7.64
$\overline{p}He^4 \to (H^3\eta_c)$	0.938	3.728	2.808	0.020	5.77	1.48	2.30
$\gamma He^4 \to (He^4\eta_c)$	0	3.728	3.728	0.120	6.59	2.24	3.96
$nHe^3 \to (He^4\eta_c)$	0.938	2.808	3.728	0.120	6.59	2.60	6.09
$dd \to (He^4\eta_c)$	1.876	1.876	3.728	0.120	6.59	2.71	9.51

Table II

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

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3. See, for example, T. Appelquist and W. Fischler, Phys. Lett. 77B, 405 (1978).
4. Due to the vector-like gluonic nature of the QCD van der Waals interaction, the pomeron scattering amplitude can be extrapolated to small s yielding a nuclear potential which incorporates multiple-gluon exchange. In principle, we could use such a procedure to evaluate the isospin-zero vector component of the low energy nucleon-nucleon potential. However, this extrapolation is not completely unambiguous if quark interchange is the dominant component, since multiple gluon exchange is difficult to distinguish from effective ω exchange. Nevertheless, in principle, the QCD van der Waals interaction provides an attractive vector-like isospin-zero potential which should be added to the usual meson-exchange potential, and this may have implications for low energy nuclear physics studies, such as nucleon-nucleon scattering and binding.
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11. See, for example, S. A. Williams *et al.*, Phys. Rev. Lett. 49, 771 (1982).

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

Figure Captions

Fig. 1. Formation of the $(c\bar{c}) - He^3$ bound state in the process $pd \to He^3X$.

Fig. 2. Decay of the $J/\psi - He^3$ bound state into $\eta_c pd$.

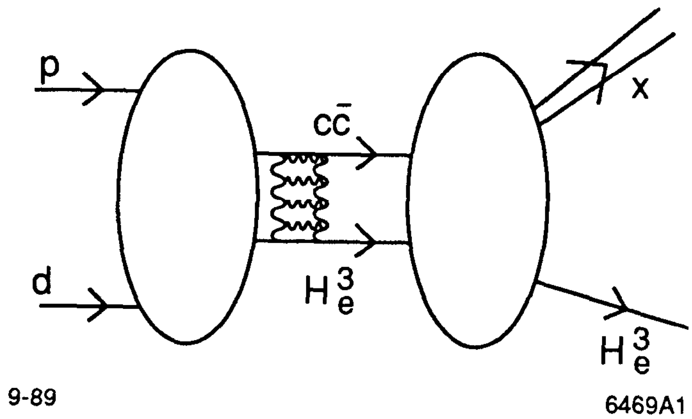

Fig. 1

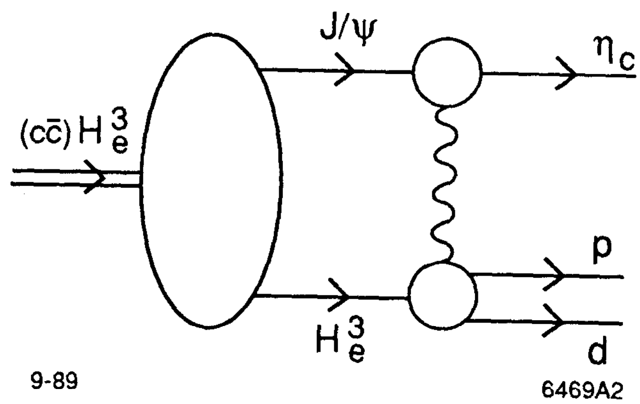

Fig. 2