Multi-quark compound states and the ^{3}He charge form factor

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The effect of multi-quark compound states on the ^{3}He charge form factor is investigated based on the Reid soft-core potential model and the relativistic harmonic oscillator quark model. A reasonable agreement with the experimental data is obtained for the momentum transfer up to ~ 1.7 GeV/c with the 2.7% six-quark compound state confined within a radius of ~ 0.9 fm.

I. INTRODUCTION

For energies up to ~ 1 GeV, most of nuclear processes and properties can be successfully described by nuclear models in which nuclei are treated as systems of nonrelativistic nucleons interacting through a two-nucleon potential or as hadronic systems composed of strongly interacting mesons and baryons. At much higher excitation energies, quark degrees of freedom (QDF) are expected to play a role, although our understanding of how the QDF will manifest in nuclei is very limited at present.

A few promising cases for studying the effect of QDF in nuclei are existences of the multi-quark compound resonances (giant hidden-color resonances, etc.)[1] as in dibaryon systems[2,3] (Kamae resonance,[4] etc.), in the charge symmetry violation[5] (the super ratio for elastic pion scatterings from ^{3}He and ^{3}H,[6] etc.), and in heavier nuclei ("anomalon"[7–9]). Other promising cases are studies of nuclei using the electromagnetic probes, such as in the cases of the European Muon Collaboration (EMC) effect[10] and of the electromagnetic structure functions and form factors of light nuclei (deuteron, ^{3}He, ^{3}H, etc.). There are now some indications that the conventional hadronic description of ^{3}He may not be sufficient in explaining the electromagnetic form factors of ^{3}He at momentum transfer squared, $Q^2 > 16$ fm^{-2},[11] as summarized by a recent review article by Friar *et al.*[12]

It is hoped that the QDF in nuclei will be understood ultimately in terms of the quantum chromodynamics (QCD). Unfortunately, we do not yet know how to formulate the nonperturbative QCD even for elementary particles, let alone for deuteron, ^{3}He, and other complex nuclei, even though the theoretical formulation and the experimental validity of the perturbative QCD are now well established for phenomena involving very large excitations and momentum transfers. Therefore, we expect that the QCD-type or QCD-motivated models of nuclei, such as the model used here, are to play a useful role in understanding the QDF in nuclei for the foreseeable future.

In an attempt to make an improvement over the conventional hadronic description[12] of the ^{3}He charge form factor, Namiki, Okano, and Oshimo[13] introduced a hybrid quark-hadron model in which the contribution[14] of the Reid soft-core potential[15] is supplemented with the relativistic harmonic oscillator quark model (RHOM) of Fujimura, Kobayashi, and Namiki[16] without making separations of the interior and exterior regions of the ^{3}He wave function. More recently, Hoodbhoy and Kisslinger[17] have applied the hybrid quark-hadron model,[18] which makes explicit separations of the interior and exterior regions, to the ^{3}He charge form factor using the Malfliet-Tjon ^{3}He wave function[19] and quark-shell model wave functions.[17] In this paper, we present the results of our calculation of the ^{3}He charge form factor, $F_{\rm ch}^{^3\rm He}(Q^2)$, employing basically the same model as Hoodbhoy and Kisslinger,[17] but using the Reid soft-core potential[15] (which is more realistic than the Malfliet-Tjon potential[19]) and the relativistic harmonic oscillator quark model of Fujimura *et al.*[16] as used by Namiki *et al.*[13] Our work may be regarded as attempts to improve the previous works of Namiki *et al.*[13] and of Hoodbhoy and Kisslinger.[17] In Sec. II, we describe in detail our model of incorporating the interior multi-quark compound (MQC) state in ^{3}He. In Sec. III, we present our explicit expressions for $F_{\rm ch}^{^3\rm He}(Q^2)$ and the zeroth (charge) component of the electromagnetic currents used. The numerical results for $F_{\rm ch}^{^3\rm He}(Q^2)$ are presented and compared with the experimental data in Sec. IV. Differences between our model and others are also discussed in this section. In Sec. V, our results are summarized and concluding remarks are given.

II. INTERIOR MULTI-QUARK COMPOUND STATES IN ^{3}He

In this section, we describe our model for incorporating multi-quark compound states in the interior region of ^{3}He, and give explicit expressions for the ^{3}He charge form factor based on our model in Sec. III.

We decompose the totally antisymmetric ^{3}He wave function as

$$\Psi(\mathbf{r}^{(1)},\boldsymbol{\rho}^{(1)})=\widetilde{\phi}(r^{(1)},\rho^{(1)})+\chi(\mathbf{r}^{(1)},\boldsymbol{\rho}^{(1)}) , \tag{1}$$

where $\widetilde{\phi}(r^{(1)},\rho^{(1)})$ are the S-state components of the ^{3}He wave function and $\chi(\mathbf{r}^{(1)},\boldsymbol{\rho}^{(1)})$ represents other higher partial wave components. The S-state wave components $\widetilde{\phi}(r^{(1)},\rho^{(1)})$ are further decomposed as

$$\tilde{\phi}(r^{(1)},\rho^{(1)})=\tilde{\phi}^{6q}_{\rm int}(r^{(1)},\rho^{(1)})+\tilde{\phi}^{9q}_{\rm int}(r^{(1)},\rho^{(1)})+\tilde{\phi}_{\rm ext}(r^{(1)},\rho^{(1)}) , \tag{2}$$

with

$$\tilde{\phi}^{6q}_{\rm int}(r^{(1)},\rho^{(1)})=\phi^{6q}_{\rm int}(r^{(1)},\rho^{(1)})\Theta(r_0-r^{(1)})\Theta(r^{(2)}-r_0)\Theta(r^{(3)}-r_0) , \tag{3}$$

$$\tilde{\phi}^{9q}_{\rm int}(r^{(1)},\rho^{(1)})=\phi^{9q}_{\rm int}\Theta(r_0-r^{(1)})\Theta(r_0-r^{(2)})\Theta(r_0-r^{(3)}) , \tag{4}$$

and

$$\tilde{\phi}_{\rm ext}(r^{(1)},\rho^{(1)})=\phi_{\rm ext}(r^{(1)},\rho^{(1)})\Theta(r^{(1)}-r_0)\Theta(r^{(2)}-r_0)\Theta(r^{(3)}-r_0) , \tag{5}$$

where r_0 is a cutoff radius in the pair coordinate, Θ denotes the Heaviside unit function, and the coordinate variables are the Lovelace variables defined as $(i,j,k$, cyclic):

$$\mathbf{r}^{(i)}=\mathbf{r}_j-\mathbf{r}_k ,$$

$$\boldsymbol{\rho}^{(i)}=\frac{1}{\sqrt{3}}(\mathbf{r}_j+\mathbf{r}_k-2\mathbf{r}_i) ,$$

$$\mathbf{R}=\sqrt{2/3}(\mathbf{r}_1+\mathbf{r}_2+\mathbf{r}_3) .$$

We note that the Jacobi coordinates are given by

$$\mathbf{x}^{(i)}=\mathbf{r}_j-\mathbf{r}_k=\mathbf{r}^{(i)} ,$$

$$\mathbf{y}^{(i)}=\mathbf{r}_i-(\mathbf{r}_j+\mathbf{r}_k)/2=-\frac{\sqrt{3}}{2}\boldsymbol{\rho}^{(i)} ,$$

$$\mathbf{R}'=\tfrac{1}{3}(\mathbf{r}_i+\mathbf{r}_2+\mathbf{r}_3)=\frac{1}{\sqrt{6}}\mathbf{R} .$$

To facilitate a model calculation for $F^{^3\rm He}_{\rm ch}(Q^2)$, we assume that $\mathcal{X}$ in Eq. (1) are the same as those given by the ^{3}He wave function[14] obtained by solving the Faddeev equation[20] with the Reid soft-core potential:[15]

$$\Psi^{\rm Reid}(\mathbf{r}^{(1)},\boldsymbol{\rho}^{(1)})=\sum_\alpha \mathscr{U}^{(\alpha)}(r^{(1)},\rho^{(1)})+\mathscr{W}(\mathbf{r}^{(1)},\boldsymbol{\rho}^{(1)}) , \tag{6}$$

i.e.,

$$\mathcal{X}(\mathbf{r}^{(1)},\boldsymbol{\rho}^{(1)})=\mathscr{W}(\mathbf{r}^{(1)},\boldsymbol{\rho}^{(1)})$$

and

$$\tilde{\phi}_{\rm ext}(r^{(1)},\rho^{(1)})=\sum_\alpha \phi^{(\alpha)}_{\rm ext}(r^{(1)},\rho^{(1)}) =\sum_\alpha \mathscr{U}^{(\alpha)}(r^{(1)},\rho^{(1)}) ,$$

where α labels the components of the exterior state. The use of Eq. (6) allows us to obtain the corresponding probabilities for $\phi^{6q}_{\rm int}$, $\phi^{9q}_{\rm int}$, and $\phi^{(\alpha)}_{\rm ext}$ as (we drop the superscripts for $r^{(1)}=r$ and $\rho^{(1)}=\rho$ from now on):

$$\langle\phi^{6q}_{\rm int}|\phi^{6q}_{\rm int}\rangle=\sum_{\alpha=1,2}\int_0^{r_0}r^2dr\int_{\rho'}^{\infty}\rho^2d\rho\,|\mathscr{U}^{(\alpha)}|^2 , \tag{7}$$

$$\langle\phi^{9q}_{\rm int}|\phi^{9q}_{\rm int}\rangle=\sum_{\alpha=1,2}\left[\int_0^{r_0}r^2dr\int_0^{\rho'}\rho^2d\rho+\int_0^{r_0}\rho^2d\rho\int_{r_0}^{r'}r^2dr\right]|\mathscr{U}^{(\alpha)}|^2 , \tag{8}$$

$$\langle\phi^{(\alpha)}_{\rm ext}|\phi^{(\alpha)}_{\rm ext}\rangle=\left[\int_{r_0}^{\infty}r^2dr\int_{r_0}^{\infty}\rho^2d\rho+\int_0^{r_0}\rho^2d\rho\int_{r'}^{\infty}r^2dr\right]|\mathscr{U}^{(\alpha)}|^2 , \tag{9}$$

where r' and ρ' are given by $(r')^2=4(r_0^2-\frac{3}{4}\rho^2)$ and $(\rho')^2=\frac{4}{3}[r_0^2-(r^2/4)]$. The summation over $\alpha=1,2$ in Eqs. (7)–(9) refers to two S-state components; the $\alpha=1$ case for the pair spin $S=1$ and pair isospin $T=0$, and the $\alpha=2$ case for $S=0$ and $T=1$.[11] For our interior multi-quark compound states, $\phi^{6q}_{\rm int}$ [Eq. (3)] and $\phi^{9q}_{\rm int}$ [Eq. (4)], we use the relativistic harmonic oscillator quark model of Fujimura, Kobayashi, and Namiki,[16] with the probabilities of each state given by Eqs. (7)–(9). Using the ^{3}He wave function of Brandenberg *et al.*,[14,21] we calculate the probabilities of the states $\phi^{6q}_{\rm int}$, $\phi^{9q}_{\rm int}$, and $\phi^{(\alpha)}_{\rm ext}$ [given by Eqs. (7)–(9), respectively], as a function of r_0, and present some results in Table I. Our calculated probabilities with $r_0=0.9$ fm are smaller than the results of

TABLE I. Probabilities of six- and nine-quark interior states in ^{3}He as a function of the cutoff radius r_0 for the case of the Reid soft-core potential. The total probability for the S states is 0.892. The remaining probability of 0.108 is for other (higher) partial-wave states.

r_0 (fm)	$\lvert\phi^{6q}_{\rm int}\rvert^2$	$\lvert\phi^{9q}_{\rm int}\rvert^2$	$\lvert\phi^{(1)}_{\rm ext}\rvert^2$	$\lvert\phi^{(2)}_{\rm ext}\rvert^2$
0.0	0.0	0.0	0.451	0.441
0.7	8.6×10^{-3}	1.4×10^{-4}	0.446	0.437
0.9	2.7×10^{-2}	9.2×10^{-4}	0.435	0.429
1.1	5.7×10^{-2}	2.5×10^{-3}	0.417	0.416

Namiki *et al.*,[13] $|\phi_{\rm int}^{6q}|^2 \cong 3.8\times10^{-2}$ and $|\phi_{\rm int}^{9q}|^2 \cong 1.6\times10^{-3}$ which are obtained by fitting the ^{3}He charge form factor. We note that $r_0=0.9$ fm is very close to $1.04a_{\rm rms}$ used by Jaffe *et al.*,[22] with $a_{\rm rms}\cong 0.88$ fm (the rms charge radius of proton) for the case of the Reid-soft-core correlation function used in their analysis of the EMC effect.[10]

Since the probabilities of higher partial wave interior states are expected to be much smaller than the probabilities of the S-wave interior states ($\alpha=1$ and 2) due to the centrifugal repulsion, our neglect of the cutoff for the higher partial wave interior state is expected to be a good approximation.

III. ^{3}He CHARGE FORM FACTOR

With the interior-exterior separation of the ^{3}He wave function described in Sec. II, we can write $F_{\rm ch}^{^3\rm He}(Q^2)$ as

$$F_{\rm ch}^{^3\rm He}(Q^2)=\cos^2\theta_1[F_{\rm N}(Q^2)+F_\pi(Q^2)] +\sin^2\theta_1[\cos^2\theta_2 F_{6q-3q}(Q^2) +\sin^2\theta_2 F_{9q}(Q^2)]\ , \qquad (10)$$

where $F_{\rm N}(Q^2)$, $F_\pi(Q^2)$, $F_{6q-3q}(Q^2)$, and $F_{9q}(Q^2)$ are contributions from the exterior impulse approximation, exterior meson-exchange charge density operators, and interior six-quark and nine-quark charge density operators, respectively. The trigonometric functions in Eq. (10) are related to the probabilities:

$$\langle \tilde\phi_{\rm int}^{6q} \,|\, \tilde\phi_{\rm int}^{6q}\rangle=\sin^2\theta_1\cos^2\theta_2\ ,$$

$$\langle \tilde\phi_{\rm int}^{9q} \,|\, \tilde\phi_{\rm int}^{9q}\rangle=\sin^2\theta_1\sin^2\theta_2\ ,$$

and

$$\sum_\alpha \langle \phi_{\rm ext}^{(\alpha)} \,|\, \phi_{\rm ext}^{(\alpha)}\rangle=\cos^2\theta_1\ .$$

For the exterior impulse approximation, the single nucleon charge density operator for the ith nucleon is given by

$$\rho_i(Q^2)=\tfrac{1}{2}[G_E^S(Q^2)+\tau_z(i)G_E^V(Q^2)]e^{-i\mathbf{Q}\cdot\mathbf{r}_i}\ , \qquad (11)$$

where $G_E^S(Q^2)$ and $G_E^V(Q^2)$ are the Sachs form factors normalized as $G_E^S(0)=G_E^V(0)=1$. The final expression for $F_{\rm N}(Q^2)$ is

$$F_{\rm N}(Q^2)=\sum_T A_T^N\left[\sum_\alpha \int\int \tilde\phi_{\rm ext}^{(\alpha)*}(r,\rho)j_0\left[\frac{\rho Q}{\sqrt3}\right]\tilde\phi_{\rm ext}^{(\alpha)}(r,\rho)r^2dr\,\rho^2d\rho+\int\int \chi^*(\mathbf{r},\boldsymbol\rho)e^{-i\mathbf{Q}\cdot\boldsymbol\rho/\sqrt3}\chi(\mathbf{r},\boldsymbol\rho)d\mathbf{r}\,d\boldsymbol\rho\right]\ , \qquad (12)$$

where T is the pair isospin and A_T^N is given by $A_0^N=\frac34[G_E^S(Q^2)+G_E^V(Q^2)]$, and $A_1^N=\frac14[3G_E^S(Q^2)-G_E^V(Q^2)]$. $\tilde\phi_{\rm ext}^{(\alpha)}$ represent two S-state components ($\alpha=1$ and 2) of $\tilde\phi_{\rm ext}$ given by Eq. (5) with the identification $\tilde\phi_{\rm ext}^{(\alpha)}(r,\rho)=\mathscr{U}^{(\alpha)}(r,\rho)$ of Eq. (6).

For the meson-exchange charge density operator, we use the zeroth (charge) component of the four-vector pair current given as[23]

$$\rho(\mathbf{k}_1,\mathbf{k}_2,\mathbf{k}_1',\mathbf{k}_2')=\frac{1}{(2\pi)^3}\delta(\mathbf{k}_1+\mathbf{k}_2-\mathbf{k}_1'-\mathbf{k}_2'-\mathbf{Q})\left[\frac{g^2}{8m^3}\right] \times\left\{[G_M^S(Q^2)\boldsymbol\tau(1)\cdot\boldsymbol\tau(2)+G_M^V(Q^2)\tau_z(2)]\frac{(\boldsymbol\sigma_1\cdot\mathbf{Q})[\boldsymbol\sigma_2\cdot(\mathbf{k}_2-\mathbf{k}_2')]}{(\mathbf{k}_2-\mathbf{k}_2')^2+\mu^2}+(1\leftrightarrow2)\right\}\ , \qquad (13)$$

or in the configuration space representation

$$\rho(\mathbf{r}_1,\mathbf{r}_2,Q)=\frac{i}{4\pi}\cdot\frac{g^2}{8m^3}(1+\mu r)\frac{e^{-\mu r}}{r^2}\{[G_M^S(Q^2)\boldsymbol\tau(1)\cdot\boldsymbol\tau(2)+G_M^V(Q^2)\tau_z(2)](\boldsymbol\sigma_1\cdot\mathbf{Q})(\boldsymbol\sigma_2\cdot\mathbf{r})e^{-i\mathbf{Q}\cdot\mathbf{r}_1}-(1\leftrightarrow2)\}\ , \qquad (14)$$

with the Sach form factors normalized as $G_M^V(0)=4.7$ and $G_M^S(0)=0.88$. The πN coupling constant is taken to be $g^2/4\pi=14$, and m is the nucleon mass.

The use of the above meson-exchange operator leads to $F_\pi(Q^2)$ given below:

$$F_\pi(Q^2)=\sum_{\alpha=1,2}A_\alpha^\pi\left[-\frac{3Q}{32\pi}\right]\left[\frac{g^2}{m^3}\right]\int\int\tilde\phi_{\rm ext}^{(\alpha)}(r,\rho)(1+\mu r)\frac{e^{-\mu r}}{r^2}j_0\left[\frac{\rho Q}{2\sqrt3}\right]j_1\left[\frac{rQ}{2}\right]\tilde\phi_{\rm ext}^{(\alpha)}(r,\rho)r^2dr\,\rho^2d\rho\ , \qquad (15)$$

where A_α^π is given by $A_1^\pi=G_M^S(Q^2)$ and $A_2^\pi=G_M^S(Q^2)+\frac13 G_M^V(Q^2)$. Although we use only the exterior S-state components for $F_\pi(Q^2)$, the approximation may be reasonable, since the higher partial contributions are expected to be small.

For calculating the contribution of the interior states to $F_{\rm ch}^{^3\rm He}(Q^2)$, we use the results of the relativistic harmonic oscillator model[16] and identify $F_{6q-3q}(Q^2)$ and $F_{9q}(Q^2)$ as

$$F_{6q-3q}(Q^2)=\langle\tilde\phi_{\rm int}^{6q}\,|\,\mathscr{O}_{\rm em}^{6q}(2,3)+\mathscr{O}_{\rm em}^{\rm N}(1)\,|\,\tilde\phi_{\rm int}^{6q}\rangle \qquad (16)$$

and

$$F_{9q}(Q^2)=\langle\tilde\phi_{\rm int}^{9q}\,|\,\mathscr{O}_{\rm em}^{9q}(1,2,3)\,|\,\tilde\phi_{\rm int}^{9q}\rangle\ , \qquad (17)$$

where the operators $\mathscr{O}_{\rm em}^{6q}(2,3)$, $\mathscr{O}_{\rm em}^{9q}(1,2,3)$, and $\mathscr{O}_{\rm em}^{\rm N}(1)$ represent the zeroth (charge) components of the electromagnetic currents for six quarks (for a pair of nucleons 2 and 3), nine quarks, and a single nucleon (nucleon 1), respectively. The explicit expressions for $F_{6q-3q}(Q^2)$ and $F_{9q}(Q^2)$ in Eqs. (16) and (17) are given by Namiki *et al.*[13] as

$$F_{nq}(Q^2)=\left[1+\frac{Q^2}{2m_n^2}\right]^{-n+1} \times\exp\left[-\left[\frac{n-1}{4\alpha_n}\right]\frac{Q^2}{(1+Q^2/2m_n^2)}\right], \quad (18)$$

with $\alpha_n=n^{3/2}K$, $K=0.096$ (GeV/c)^2,

$m_3=1.097$ GeV/c, $m_6=1.2$ GeV/c ,

$m_9=1.5$ GeV/c ,

and

$$F_{6q-3q}(Q^2)=\tfrac{1}{3}[2F_{6q}(Q^2)G_{3q}(Q^2) +F_{3q}(Q^2)G_{6q}(Q^2)]F_{2c}(Q^2)\ , \quad (19)$$

where

$$G_{nq}(Q^2)=\left[1+\frac{Q^2}{2m_n^2}\right]^{-n+1}$$

and

$$F_{2c}(Q^2)=\left[1+\frac{Q^2}{m^2}\right]^{-1}$$

with $m=1$ GeV/c. We note that $F_{nq}(Q^2)$ given by Eq. (18) has an appropriate dimensional asymptotic scaling.[24] We also note that the above expression for $F_{6q-3q}(Q^2)$ given by Eq. (19) is different from the corresponding expression used by Hoodbhoy and Kisslinger [Eq. (4) of Ref. 17]. The phase in Eqs. (18) and (19) cannot be determined by the present model. As done by Namiki *et al.*,[13] we take the negative sign for both $F_{9q}(Q^2)$ and $F_{6q-3q}(Q^2)$.

IV. RESULTS

We calculate the ^{3}He charge form factor using Eqs. (10)—(19), described in Sec. III. For the electromagnetic nucleon form factors, G_E^S, G_E^V, G_M^S, and G_M^V, we use the five-parameter dipole fits of Iachello *et al.*[25] Our calculated results are shown schematically and compared with the experimental data[26,27] in Figs. 1 and 2. We note that, for $Q^2 \geq 20$ fm^{-1}, the experimental data[27] are for $|A^{1/2}(Q^2)|$, which is defined as

$$A(Q^2)=(|F_{\rm ch}^{^3{\rm He}}(Q^2)|^2+\mu_{^3{\rm He}}^2\tau|F_{\rm mag}^{^3{\rm He}}(Q^2)|^2)/(1+\tau)$$

with $\mu_{^3{\rm He}}=-3.2$ nuclear magnetons and $\tau=Q^2/(4M^2)$ (M is ^{3}He mass), rather than for $F_{\rm ch}^{^3{\rm He}}(Q^2)$. Since there are no available data for $F_{\rm ch}^{^3{\rm He}}(Q^2)$ [or $F_{\rm mag}^{^3{\rm He}}(Q^2)$], $Q>6$ fm^{-1},[11] comparison of the calculated $F_{\rm ch}^{^3{\rm He}}(Q^2)$ and the experimental $A^{1/2}(Q^2)$ may be premature, but we assume that $A(Q^2)\approx|F_{\rm ch}^{^3{\rm He}}(Q^2)|^2$ for 10 fm$^{-1}>Q\gtrsim 6$ fm^{-1}.

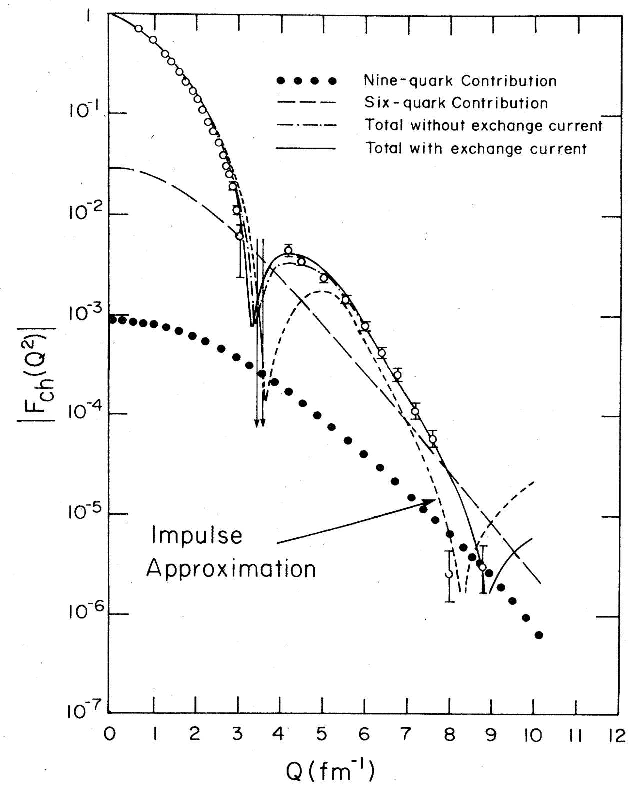

FIG. 1. Comparison of the calculated results and the experimental data (Refs. 26 and 27, open circles) for the ^{3}He charge form factor, for the case of the cutoff confinement radius of $r_0=0.9$ fm. Contributions from the exterior three-quark (impulse approximation, short dash), interior six-quark (long dash), and nine-quark (dots) states are plotted individually, while total contributions are indicated by a dash-dot curve (without exchange current) and a solid curve (with exchange current).

The above discussion shows the importance of measuring $F_{\rm mag}^{^3{\rm He}}(Q^2)$, $Q\gtrsim 6$ fm^{-1} in the future.

In Fig. 1, we show the effect of individual contributions to the ^{3}He charge form factor for the case of $r_0=0.9$ fm. As is well known,[14] the calculated result of the impulse approximation, $F_{\rm N}(Q^2)$, does not agree with the experimental data for the momentum transfer, $Q\gtrsim 3$ fm^{-1}. Our calculated result of $F_{\rm ch}^{^3{\rm He}}(Q^2)$ with the addition of the six-quark contribution, $F_{6q-3q}(Q^2)$, to the impulse approximation, $F_{\rm N}(Q^2)$, improves substantially the agreement with the experimental data. The meson-exchange contribution, $F_\pi(Q^2)$, is rather small, as expected from the fact that the cutoff of the meson-exchange operators at $r_0=0.9$ fm eliminate the substantial part of these operators which have the pion range of 1.4 fm. The same agreement can be used to justify, to a good approximation, our exclusion of heavier meson-exchange contributions, two-pion exchange contribution, and the effect of three-nucleon forces,[28] etc. Even for the meson-exchange contributions, we can avoid the ambiguity associated with the cutoff parameter of the πNN form factor for the exchange current. The widely used monopole form of the πNN form factor has a cutoff parameter Λ ranging from

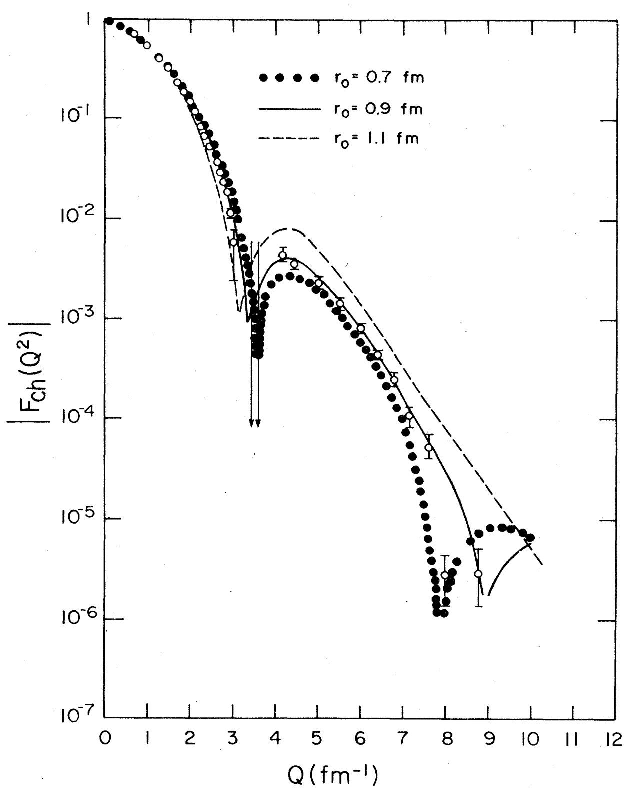

FIG. 2. Comparison of the calculated results of the ^{3}He charge form factor with $r_0=0.7$ fm (dots), 0.9 fm (solid curve), and 1.1 fm (dashed curve). Open circles are for the experimental data.

4μ to 8.5μ (μ is the pion mass).[11] Our results are insensitive to Λ.

The nine-quark contribution $F_{9q}(Q^2)$ to $F_{ch}^{^3\text{He}}(Q^2)$ is small even for $Q\cong 6$–10 fm^{-1}, which is consistent with the fact that $F_{ch}^{^3\text{He}}(Q^2)$ for $Q\cong 6$–10 fm^{-1} does not exhibit dimensional asymptotic scaling as a nine-quark system.[24]

To study the dependence of $F_{ch}^{^3\text{He}}(Q^2)$ on r_0, we calculate $F_{ch}^{^3\text{He}}(Q^2)$ with $r_0=0.7$, 0.9, and 1.1 fm and compare the calculated results with each other and the experimental data. As we increase r_0, the probabilities of six- and nine-quark compound states increase (see Table I) and both contributions to $F_{ch}^{^3\text{He}}(Q^2)$ become enhanced. However, the six-quark contribution is still dominant over the nine-quark contribution for $Q \lesssim 10$ fm^{-1} considered in this work. The effect of the meson-exchange contributions to $F_{ch}^{^3\text{He}}(Q^2)$ becomes smaller as r_0 increases. Assuming that $A^{\exp}(Q^2)\approx |F_{ch}^{^3\text{He}}(Q^2)|^2$, for $Q \gtrsim 6$ fm^{-1}, the calculated $F_{ch}^{^3\text{He}}(Q^2)$ with $r_0=0.9$ fm seems to give a better fit to the experimental data than the results with $r_0=0.7$ or 1.1 fm. As noted before, $r_0=0.9$ fm is close to the six-quark bag radius used by Jaffe *et al.*[22] in their analysis of the EMC effect.[10]

There have been previous calculations[13,19,29] of $F_{ch}^{^3\text{He}}(Q^2)$ similar to ours. Namiki *et al.*[13] used a similar model without introducing the confinement cutoff radius, r_0, and without the meson-exchange pair current. The meson-exchange effect is small and hence can be neglected to a good approximation, but the neglect of the confinement radius for six-quark and nine-quark states cannot be justified on physical grounds. Our results with $r_0=0.9$ fm can be regarded as an improvement over their calculations and give a better fit to the experimental data for $F_{ch}^{^3\text{He}}(Q^2)$. We note that we determine the probabilities of the interior quark states as a function of r_0, while they obtain them by fitting their calculated $F_{ch}^{^3\text{He}}(Q^2)$ to the experimental data.

Other similar calculations of $F_{ch}^{^3\text{He}}(Q^2)$ are those of Refs. 17 and 29. Although their results for $F_{ch}^{^3\text{He}}(Q^2)$ with $r_0=1$ fm give a quantitatively similar fit as ours in the range of $Q^2<20$ cm^{-2}, the reasons for the agreement are different. They use the ^{3}He wave function generated from the Malfliet-Tjon potential,[19] while we use the Reid soft-core potential.[15] Therefore, their result for the impulse approximation is quite different from ours for $Q^2>10$ fm^{-2}, which may explain why they need a substantially larger six-quark probability (12%–15% for $r_0=1$ fm) and contribution, while we need only a 2.7% six-quark state. A larger six-quark probability for the Malfliet-Tjon potential model[19] is expected since it has a much weaker short-range repulsion than the Reid soft-core potential.[15]

V. SUMMARY AND CONCLUSION

With the bound-state ^{3}He wave function, which has the interior multi-quark state confined within a cutoff radius r_0 and the exterior three-nucleon state, we have calculated the ^{3}He charge form factor using the RHOM for the interior state and a modified Reid soft-core ^{3}He wave function for the exterior state (with cutoff S states). The probabilities of the interior six-quark and nine-quark states are determined from the missing part of the original Reid soft-core ^{3}He wave function. We find a reasonable fit of our calculated $F_{ch}^{^3\text{He}}(Q^2)$ to the experimental data for $Q\lesssim 10$ fm^{-1}, with $r_0=0.9$ fm and with small probabilities of the six-quark (2.7%) and nine-quark (0.03%) interior states. Our results of a small nine-quark contribution to $F_{ch}^{^3\text{He}}(Q^2)$ for $Q\lesssim 10$ fm^{-1} (or $Q\lesssim 2$ GeV/c) are consistent with the fact that the experimental data for $F_{ch}^{^3\text{He}}(Q^2)$ have not reached the dimensional scaling asymptotic region,[24,27] and corresponds to a situation somewhere between two extreme views[30,31] of the effect of the perturbative quantum chromodynamics for the elastic form factors at these momentum transfers.

The separation of the ^{3}He wave function into the interior and exterior parts can remove many ambiguities associated with the short-range correlations, such as ambiguity of the πNN form factor in the nuclear forces and meson-exchange currents. It can also suppress heavy-meson exchange and two-pion exchange effects, as well as the three-nucleon force effect, in favor of the confined multi-quark currents and dynamics. Unfortunately, we do not yet have a more consistent model of the three-nucleon systems which includes explicit quark degrees of freedom, but may be able to develop a more sophisticated description based on the resonating group quark-cluster model of nuclei[1–3] in the near future.

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