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Pion Contributions to the Electromagnetic Interaction with a Nucleus

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A theoretical approach to describing the electromagnetic interaction in hadronic field theories containing charged mesons at finite baryon density is presented. Path-integral methods and bilocal auxiliary fields are used to identify covariant, gauge-invariant levels of approximation, which sum infinite classes of Feynman diagrams. At the lowest level an isolated nucleon acquires a nontrivial electromagnetic structure, and the finite-density nucleus is described as a noninteracting system of quasi particles with that structure.

Many descriptions of the electromagnetic (EM) interaction with a nucleus center around the approximate treatment of a nucleus as a weakly or noninteracting system of quasi particles. The aim of this work is to initiate the development of a consistent framework within which the EM properties of a noninteracting system of dressed nucleons can be investigated, and to which corrections above the noninteracting level can be systematically made. The approach taken here is to employ bilocal auxiliary fields and path-integral methods to formulate a gauge invariant expansion in fluctuations above the vacuum, which are then allowed to interact with the system of valence nucleons. Previous work in a quark-meson framework demonstrates how the fluctuations of $q\bar{q}$ bilocal fields about their vacuum values are related to mesons[1] and, through the incorporation of a set of chemical potentials, how a system of valence quarks interacting with these mesons can be described[2]. Similar techniques are applied here in a purely hadronic setting to describe density-dependent collective excitations[3] which are expected to provide important contributions to EM processes beyond the level of a Fermi gas. The auxiliary field approach employed here should therefore be viewed as an organizational tool for systematically incorporating the effects of these collective modes in an expansion above the noninteracting Fermi gas of quasi particles dressed by the vacuum, which arises at zeroth order, and not as a representation of physical particles. Chemical potentials will again be employed to fix the number of valence nucleons and to separate their contribution to physical quantities from those arising from the vacuum.

The calculation of physical quantities is achieved here through the use of the partition function for the system, Z, written in the path-integral formulation. Derivatives of the natural logarithm of the partition function with respect to source fields give expectation values of the associated currents. In the presence of nonzero chemical potentials this procedure results in expectation values of currents for a system of valence nucleons. The incorporation of a background EM gauge field provides a source for the EM current. Expectation values of the EM current are thereby obtained. The gauge invariance of the partition function ensures that the physical quantities obtained in this way are gauge invariant and conserved.

In purely hadronic models such as quantum hadrodynamics[4] (QHD) it has been demonstrated that the charged pions play a significant role in the description of the EM form factors of the nucleon[5], and are expected to make important contributions to the EM properties of nuclear systems. The radiative pion dressing of the nucleon presents complications because of the isovector nature of the pion, and therefore requires a more general approach to include the effects of charged particle exchange than is needed for uncharged mesons. In this initial

investigation, we will therefore concentrate on the incorporation of charged pions into a description of the EM properties of a system of nucleons in a hadronic field theory. It is expected that the effects of other mesons can be included in a similar manner. The method of calculating EM properties based on functional differentiation of the partition function will be demonstrated at the level of a Fermi gas of nucleons having radiative pion dressing.

We begin with an action describing the interaction of charged pions, ϕ and $\phi^\dagger$, with nucleons, $\psi^t = (p, n)$, given in the Euclidean metric by

$$S\left[\bar{\psi},\psi,\phi^\dagger,\phi\right] \equiv \int d^4x \left\{ \bar{\psi}\left[\not{\partial}+m+ig\gamma_5\left(\tau\phi^\dagger+\tau^\dagger\phi\right)\right]\psi+\left(\partial_\mu\phi^\dagger\right)\left(\partial_\mu\phi\right)+m_\phi^2\phi^\dagger\phi \right\}, \tag{1}$$

where $\tau = (\tau_1 - i\tau_2)/\sqrt{2}$ and τ_1 and τ_2 are the usual Pauli spin matrices. Clearly a more realistic description would have to include other meson fields such as those required to impose chiral symmetry, and baryons such as the delta. However, for the sake of illustrating this method of treating dressing due to a charged particle, only the charged pion and nucleon fields are used here.

To study the EM interaction, a background gauge field can be incorporated by the minimal substitution prescription $\partial_\nu \to \partial_\nu - iqA_\nu$, where q is the relevant charge. The gauge invariant action is then given by

$$S\left[\bar{\psi},\psi,\phi^\dagger,\phi,A_\nu\right] = \left(\bar{\psi},G^{-1}[A_\nu]\psi\right) + \left(\phi^\dagger,\Delta^{-1}[A_\nu]\phi\right) + \left(\phi^\dagger,j\right) + \left(j^\dagger,\phi\right), \tag{2}$$

where the notation (A, B) implies $\int d^4x A(x)B(x)$, and where

$$\begin{aligned} G^{-1}[A_\nu] &\equiv \not{\partial} + m - i\hat{Q}\not{A} \\ \Delta^{-1}[A_\nu] &\equiv \left(\overleftarrow{\partial}_\nu + iA_\nu\right)\left(\overrightarrow{\partial}_\nu - iA_\nu\right) + m_\phi^2 \\ j^\dagger &\equiv ig\bar{\psi}\gamma_5\tau^\dagger\psi. \end{aligned} \tag{3}$$

The nucleon charge matrix $\hat{Q}$ is defined by $\hat{Q} \equiv (1+\tau_3)/2$. The partition function (Euclidean space generating functional) is given by

$$Z[A_\nu] = \int D\bar{\psi}D\psi D\phi^\dagger D\phi e^{-S[\bar{\psi},\psi,\phi^\dagger,\phi,A_\nu]}, \tag{4}$$

and is trivially gauge invariant. The motivation in what follows is to eliminate the nucleon and pion fields in favor of auxiliary fields characterizing collective degrees of freedom.

Upon performing the integration over the pion fields, $\phi^\dagger$ and ϕ, the action exponent S becomes

$$S\left[\bar{\psi},\psi,A_\nu\right] = \left(\bar{\psi},G^{-1}[A_\nu]\psi\right) - \left(j^\dagger,\Delta[A_\nu]j\right) + Tr Ln \Delta^{-1}[A_\nu]. \tag{5}$$

The $Tr Ln\Delta^{-1}[A_\nu]$ term corresponds to closed pion loops, and can therefore only contribute to the EM properties of the vacuum. In the spirit of considering the properties of a system of valence nucleons, all terms in which pions are not connected to nucleons will be neglected. Hence, the last term in (5) is dropped. The second term in (5) can be rewritten using the Fierz reordering $(\gamma_5)_{ij}(\gamma_5)_{kl} = \Lambda^\theta_{il}\Lambda^\theta_{kj}$, where $\Lambda^\theta = \frac{1}{2}\{1,\gamma_5,\gamma_\nu\gamma_5,i\gamma_\nu,\sigma_{\nu\eta}\}$, to obtain

$$-j^\dagger(x)\Delta(x,y;A_\nu]j(y) = -2g^2\bar{p}(x)\Lambda^\theta p(y)\Delta(x,y;A_\nu]\bar{n}(y)\Lambda^\theta n(x). \tag{6}$$

At this point we must depart from the previous work since the two bilocal currents $\bar{p}(x)\Lambda^\theta p(y)$ and $\bar{n}(y)\Lambda^\theta n(x)$ are not equivalent, and hence the standard bosonization[6] can not be invoked. We can however utilize complex bilocal Bose fields to achieve the bosonization by inserting the resolution of unity

$$\begin{aligned}1 = \mathcal{N}\int DB^* DB \exp\Bigg[-\int d^4x d^4y &\left(B^\theta(x,y)^* + g\left[2\Delta(x,y;A_\nu]\right]^{1/2}\bar{p}(x)\Lambda^\theta p(y)\right)\\ &\times\left(B^\theta(x,y) + g\left[2\Delta(x,y;A_\nu]\right]^{1/2}\bar{n}(y)\Lambda^\theta n(x)\right)\Bigg]\end{aligned} \tag{7}$$

into the integrand of the partition function in (4), where $DB = \prod_\theta DB^\theta$. The partition function becomes

$$Z[A_\nu] = \mathcal{N}\int D\bar{p}DpD\bar{n}DnDB^*DB\, e^{-S}, \tag{8}$$

where the action exponent $S = S[\bar{p},p,\bar{n},n,B^*,B,A_\nu]$ is given by

$$\begin{aligned}S = \int d^4x d^4y\ \Big\{&\bar{p}(x)G_p^{-1}(x,y;A_\nu]p(y) + \bar{n}(x)G_n^{-1}(x,y;A_\nu]n(y)\\ &+ B^\theta(x,y)^* B^\theta(x,y)\Big\}.\end{aligned} \tag{9}$$

The inverse propagators appearing in (9) are defined as

$$G_p^{-1}(x,y;A_\nu] \equiv \left(\not{\partial} + m - i\not{A}\right)\delta(x-y) + g\left(2\Delta(x,y;A_\nu]\right)^{1/2}\Lambda^\theta B^\theta(x,y) \tag{10}$$

and

$$G_n^{-1}(x,y;A_\nu] \equiv \left(\not{\partial} + m\right)\delta(x-y) + g\left(2\Delta(y,x;A_\nu]\right)^{1/2}\Lambda^\theta B^\theta(y,x)^*. \tag{11}$$

To fix the number of *physical* (dressed) valence nucleons and to isolate their contribution to the partition function one can[2] use a canonical transformation to incorporate chemical potentials, μ_p and μ_n for the protons and neutrons respectively. Then integration over the Grassmann fields with the appropriate adjustment of boundary condition produces the grand partition function $Z[A_\nu,\mu_p,\mu_n]$ given by

$$Z[A_\nu,\mu_p,\mu_n] = \mathcal{N}\int DB^* DB e^{-S[B^*,B,A_\nu,\mu_p,\mu_n]}, \tag{12}$$

where the bosonized action exponent $S[B^*,B,A_\nu,\mu_p,\mu_n] = S_{val} + S_{vac}$ is given by

$$S_{val} = -Tr\left[\left(LnG_p^{-1}[\mu_p] - LnG_p^{-1}[0]\right) + \left(p\rightarrow n\right)\right], \tag{13}$$

$$S_{vac} = -Tr\left(LnG_p^{-1}[0] + LnG_n^{-1}[0]\right) + \int d^4x d^4y B^\theta(x,y)^* B^\theta(x,y). \tag{14}$$

The separation in Eqs.(13) and (14) isolates the valence nucleon contributions from the vacuum action. The inverse propagators appearing in (13) and (14) are defined[2] by $G_{p,n}^{-1}(x,y;A_\nu,\mu_{p,n}] \equiv e^{\mu_{p,n}x_4}G_{p,n}^{-1}(x,y;A_\nu]e^{-\mu_{p,n}y_4}$ with $G_{p,n}^{-1}$ given in (10) and (11).

The only remaining integration in the partition function is over the bilocal auxiliary fields. For modest energies these fields can be evaluated at their vacuum values, B_0^* and B_0, corresponding to the saddle point configurations of the vacuum action, S_{vac} in (14). Corrections due to vacuum fluctuations can be considered by expanding the vacuum action about the saddle-point configurations as

$$S_{vac} = S_{vac}[B_0^*,B_0] + \int \hat{B}^*\frac{\delta S_{vac}}{\delta B_0^*\delta B_0}\hat{B} + \cdots, \tag{15}$$

where $\hat{B}^*$ and $\hat{B}$ are the fluctuations away from the saddle point, and the schematic notation used in (15) implies that all types of the bilocal fields (labeled by θ) are to be included. In the analogous bosonization of a quark field theory these fluctuations, when treated at the mean field level, result in non-local potentials[2]. Without the source provided by the valence nucleons these fluctuations would vanish by construction. Therefore, in regions of low source density one would expect these fluctuations to be small. The result of neglecting these fluctuations is that the system is described as a gas of noninteracting quasi particles having radiative dressing from the vacuum as we now show.

The saddle-point configurations of the vacuum action are defined by $\delta S_{vac}/\delta B_0^{\#} = 0$ and $\delta S_{vac}/\delta B_0^{*\#} = 0$. These are implicitly dependent on the gauge field, A_ν. In terms of the saddle-point configurations the gauge-field-dependent self energies for the nucleons are defined as

$$
\begin{aligned}
\Sigma_p(x,y;A_\nu] &\equiv g\Big(2\Delta(x,y;A_\nu]\Big)^{1/2}\Lambda^{\#}B_0^{\#}(x,y;A_\nu] \\
&= 2g^2\Delta(x,y;A_\nu]\gamma_5\mathcal{G}_n(x,y;A_\nu]\gamma_5
\end{aligned}
\tag{16}
$$

and

$$
\begin{aligned}
\Sigma_n(x,y;A_\nu] &\equiv g\Big(2\Delta(y,x;A_\nu]\Big)^{1/2}\Lambda^{\#}B_0^{*\#}(y,x;A_\nu] \\
&= 2g^2\Delta(y,x;A_\nu]\gamma_5\mathcal{G}_p(x,y;A_\nu]\gamma_5,
\end{aligned}
\tag{17}
$$

where $\mathcal{G}_p$ and $\mathcal{G}_n$ self consistently contain Σ_p and Σ_n respectively. Eqs.(16) and (17) are the Schwinger-Dyson forms for the nucleon self energy with bare nucleon-pion vertices. With the bilocal fields evaluated at the saddle-point level, the partition function can be written as $Z^{(0)}[A_\nu,\mu_p,\mu_n] = \exp(W^{(0)}[A_\nu,\mu_p,\mu_n])$ where

$$
W^{(0)}[A_\nu,\mu_p,\mu_n] \equiv -\Big[S_{val} + S_{vac}\Big]_{B_0^*,B_0}
\tag{18}
$$

and where the inverse propagators occurring in the action at this level are given by

$$
\begin{aligned}
G_p^{-1}(x,y;A_\nu,\mu_p] &= e^{\mu_p x_4}\mathcal{G}_p^{-1}(x,y;A_\nu]e^{-\mu_p y_4} \\
&= \Big(\not\partial + m - \gamma_4\mu_p - i\not A\Big)\delta(x-y) + e^{\mu_p x_4}\Sigma_p(x,y;A_\nu]e^{-\mu_p y_4}
\end{aligned}
\tag{19}
$$

and

$$
\begin{aligned}
G_n^{-1}(x,y;A_\nu,\mu_n] &= e^{\mu_n x_4}\mathcal{G}_n^{-1}(x,y;A_\nu]e^{-\mu_n y_4} \\
&= \Big(\not\partial + m - \gamma_4\mu_n\Big)\delta(x-y) + e^{\mu_n x_4}\Sigma_n(x,y;A_\nu]e^{-\mu_n y_4}.
\end{aligned}
\tag{20}
$$

Expectation values of the currents coupled to A_ν, μ_p, and μ_n can now be obtained by differentiating the functional $W^{(0)}[A_\nu,\mu_p,\mu_n]$. To obtain finite quantities, one must subtract the corresponding vacuum expectation values defined by taking $\mu_p = \mu_n = 0$. For example, the finite values of the nucleon occupation numbers, N_p and N_n, are given by

$$
\begin{aligned}
N_{p,n} &\equiv \frac{1}{\beta}\left[\frac{\partial W^{(0)}}{\partial\mu_{p,n}} - \frac{\partial W^{(0)}}{\partial\mu_{p,n}}\bigg|_{\mu_{p,n}=0}\right]_{A_\nu=0} \\
&= \frac{1}{\beta}Tr\left[G_{p,n}[\mu_{p,n}]\frac{\partial G_{p,n}^{-1}[\mu_{p,n}]}{\partial\mu_{p,n}} - G_{p,n}[0]\,\frac{\partial G_{p,n}^{-1}[\mu_{p,n}]}{\partial\mu_{p,n}}\bigg|_{\mu_{p,n}=0}\right]_{A_\nu=0},
\end{aligned}
\tag{21}
$$

where $\beta = \int dx_4$. At the saddle-point level of approximation, and with the gauge field $A_\nu = 0$, the propagators are translationally invariant allowing a spectral representation to be employed[2]. The occupation numbers given in (21) can thereby be evaluated to give

$$
N_{p,n} = V\int\frac{d^3p}{(2\pi)^3}\Theta\Big(\mu_{p,n} - \epsilon_{p,n}^{(+)}(\vec p)\Big),
\tag{22}
$$

where $\epsilon_{p,n}^{(+)}(\vec p)$ are the positive eigenenergies.

A quantity of particular interest for the investigation of inelastic electron scattering is the polarization tensor, which is defined here as

$$
\Pi_{\alpha\beta}(x,y) \equiv \frac{\delta W^{(0)}[A_\nu,\mu_p,\mu_n]}{\delta A_\alpha(x)\delta A_\beta(y)}\bigg|_{A_\nu=0} - \frac{\delta W^{(0)}[A_\nu,0,0]}{\delta A_\alpha(x)\delta A_\beta(y)}\bigg|_{A_\nu=0}.
\tag{23}
$$

The second term in (23) is clearly the polarization of the vacuum. For the sake of illustration, we evaluate this quantity for the case of zero occupied neutron levels, $\mu_n = 0$, and simply take $\mu_p = \mu$. We also consider only the first order dependence of the inverse propagators in $W^{(0)}$ on the gauge field, A_ν. The translational invariance of the saddle-point approximation then allows us to write $\Pi_{\alpha\beta}(x,y) \equiv \Pi_{\alpha\beta}^{(\mu)}(x-y) - \Pi_{\alpha\beta}^{(0)}(x-y)$ where

$$
\begin{aligned}
\Pi_{\alpha\beta}^{(\mu)}(r) = &-\int_{-\infty+i\mu}^{\infty+i\mu}\frac{dp_4dk_4}{(2\pi)^2}\int\frac{d^3pd^3k}{(2\pi)^6}e^{-i(p-k)r} \\
&\times\mathrm{tr}\left[G_0(p)\Gamma_\alpha^{(p)}\Big(\frac{p+k}{2},p-k\Big)G_0(k)\Gamma_\beta^{(p)}\Big(\frac{p+k}{2},k-p\Big)\right]
\end{aligned}
\tag{24}
$$

and where the photon-proton vertex, $\Gamma_\alpha^{(p)}(P,q)$ is defined by

$$\left.\frac{\delta \mathcal{G}_p^{-1}(p,k;A_\nu]}{\delta A_\alpha(q)}\right|_{A_\nu=0} \equiv \frac{1}{(2\pi)^2}\delta(p-k-q)\Gamma_\alpha^{(p)}\left(\frac{p+k}{2},q\right), \tag{25}$$

and $\mathcal{G}_p(p,k;A_\nu]_{A_\nu=0} \equiv \delta(p-k)G_0(p)$. The inverse propagator appearing in (25) is obtained from (19) with $\mu_p = 0$. It is evident, through examination of (19) and (16), that the photon-proton and photon-neutron vertex functions satisfy coupled integral equations.

By employing the spectral representation of the propagators one obtains from (24)

$$\Pi_{\alpha\beta}^{(\mu)}(r) = \int \frac{d^3p d^3k}{(2\pi)^6} e^{-i(\vec{p}-\vec{k})\cdot\vec{r}} \int_{-\infty+i\mu}^{\infty+i\mu} \frac{dp_4}{2\pi i}\frac{e^{-ip_4r_4}}{p_4 - i\epsilon(\vec{p})} \times \int_{-\infty+i\mu}^{\infty+i\mu} \frac{dk_4}{2\pi i}\frac{e^{ik_4r_4}}{k_4 - i\epsilon(\vec{k})}\Omega_{\alpha\beta}(p;k), \tag{26}$$

where $\Omega_{\alpha\beta}(p;k) \equiv -Z^{-2}\left[\bar{u}(\vec{p})\Gamma_\alpha^{(p)}\left(\frac{p+k}{2},p-k\right)u(\vec{k})\right]\left[\bar{u}(\vec{k})\Gamma_\beta^{(p)}\left(\frac{p+k}{2},k-p\right)u(\vec{p})\right]$. The wave function renormalization, Z results from the frequency dependence of the self energy. The wave-function renormalization constant is also present in the vertex function $\Gamma_\alpha^{(p)}$, defined in (25), and will be canceled in $\Omega_{\alpha\beta}$. This cancelation is a consequence of gauge invariance[7], and reflects the consistency of this approach. With the fourth component of the momenta fixed by the frequency integrations in (26), the spinors in $\Omega_{\alpha\beta}$ are solutions to the free Dirac equation with a renormalized mass arising from the vacuum self energy provided by the pion. The translational invariance of the saddle-point (vacuum) level of approximation then allows the current, $j_\alpha(\vec{p},\vec{k}) \equiv -Z\bar{u}(\vec{p})\Gamma_\alpha^{(p)}\left(\frac{p+k}{2},p-k\right)u(\vec{k})$ to be written in the usual form

$$j_\alpha(\vec{p},\vec{k}) = \bar{u}(\vec{p})\left[i\gamma_\alpha F_1(Q^2) + i\sigma_{\alpha\beta}Q_\beta F_2(Q^2)\right]u(\vec{k}), \tag{27}$$

where $Q_\alpha = p_\alpha - k_\alpha$, $p_4 = i\epsilon(\vec{p})$ and $k_4 = i\epsilon(\vec{k})$.

By performing the frequency integrations in (26) and transforming to momentum space, one obtains

$$\Pi_{\alpha\beta}^{(\mu)}(\vec{q},\omega) = \int \frac{d^3p}{(2\pi)^3}\Omega_{\alpha\beta}\left(\vec{p},i\epsilon(\vec{p});\vec{p}+\vec{q},i\epsilon(\vec{p}+\vec{q})\right) \times \frac{\Theta\left(\mu-\epsilon(\vec{p})\right)\Theta\left(\epsilon(\vec{p}+\vec{q})-\mu\right)-\Theta\left(\epsilon(\vec{p})-\mu\right)\Theta\left(\mu-\epsilon(\vec{p}+\vec{q})\right)}{\omega-\epsilon(\vec{p}+\vec{q})+\epsilon(\vec{p})-i\delta}, \tag{28}$$

where we have made the continuation $q_4 \to i\omega + \delta$. Finally, subtracting the contribution for $\mu = 0$ and taking $\epsilon(\vec{p}+\vec{q}) > \mu$, which is the relevant condition for quasi-elastic scattering, we find

$$\Pi_{\alpha\beta}(\vec{q},\omega) = -\int \frac{d^3p}{(2\pi)^3} j_\alpha(\vec{p},\vec{p}+\vec{q}) j_\beta(\vec{p}+\vec{q},\vec{p}) \frac{\Theta\left(\mu-\epsilon^{(+)}(\vec{p})\right)\Theta\left(\epsilon^{(+)}(\vec{p}+\vec{q})-\mu\right)}{\omega-\epsilon^{(+)}(\vec{p}+\vec{q})+\epsilon^{(+)}(\vec{p})-i\delta}. \tag{29}$$

This is in fact the result which is obtained in the Fermi gas model[8] with the free nucleon form factors used in the current. Here this result occurs as the lowest level in an expansion to which systematic corrections can be made.

In summary, an initial investigation of a systematic, gauge-invariant approach to considering the EM interaction with a nucleus has been presented. The approach centers around a transformation from nucleon and meson degrees of freedom to collective degrees of freedom. The fluctuations of these about their translationally invariant vacuum values are expected to be small, but give the leading corrections to the description of the nucleus as a Fermi gas of quasi particles having radiative dressing from the vacuum. The principle result of this initial investigation is summarized by Eqs.(27) and (29), and the statement that charged-pion contributions to the EM properties of the nucleus can be consistently included through this approach. Further work, including the incorporation of uncharged mesons, is being pursued and a detailed article is planned.

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