Zero Modes in Electromagnetic Form Factors of the Nucleon in a Light-Cone Diquark Model

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Abstract

We use a diquark model of the nucleon to calculate the electromagnetic form factors of the nucleon described as a scalar and axialvector diquark bound state. We provide an analysis of the zero-mode contribution in the diquark model. We find there are zero-mode contributions to the form factors arising from the instantaneous part of the quark propagator, which cannot be neglected compared with the valence contribution but can be removed by the choice of wave function. We also find that the charge and magnetic radii and magnetic moment of the proton can be reproduced, while the magnetic moment of the neutron is too small. The dipole shape of the form factors, $G_M^p(Q^2)/\mu_p$ and $G_M^n(Q^2)/\mu_n$, can be reproduced. The ratio $\mu G_E^p/G_M^p$ decreases with $Q^2$, but too fast.
1 Introduction

Methods using relativistic formalisms directly without Foldy-Wouthuysen reduction have been suggested [1, 2, 3, 4] to include relativistic effects to all orders, while avoiding truncated power expansions in $v/c$ or $p/m$. Among these relativistic formalisms, light-cone quark models are widely used. There are several works where the electroweak properties of mesons have been studied using light-cone quark models [5, 6, 7, 8, 9, 10]. Applications of the Bakamjian-Thomas construction of the relativistic few-body system [3, 4, 11, 12] to extend nonrelativistic wave functions to the relativistic domain are close to the naive quark model. Weber et al. [13, 14] give relativistic spin-flavor wave functions of a Lorentz covariant quark model. Araujo et. al. [15] connect these wave functions with the triangle Feynman diagram and give an effective Lagrangian of the nucleon, which corresponds to a field-theoretical formulation of the form factors.

From the Eqs. (9,10) in Ref. [13] (KW, for short. See also Eq. (3) below.), we see that the nucleon wave function is split into two parts as $(qq)q$. At this point, the diquark model can be introduced naturally. The diquark model has been widely used to study the property of baryons. For example, nucleon properties are studied through the Bethe-Salpeter equation with scalar and axialvector diquark correlations [16]. Kroll et. al. study the electromagnetic form factors of the proton in the time-like region and two-photon annihilation into proton-antiproton in the infinite momentum frame [17].

In all investigations of baryon properties [3, 4, 11, 12, 13, 14, 15], zero modes have not been considered. The zero-mode contributions can be interpreted as residues of virtual pair creation processes in the $q^+ (= q^0 + q^3) \rightarrow 0$ limit of the virtual photon momentum [18, 19]. In the absence of zero-mode contributions, the hadron form factor can be obtained in a straightforward way by taking into account only the valence quark contributions. In an effort to clarify this issue, Jaus [20] and Choi et. al. [21] independently investigated the electroweak form factors of spin-1 mesons in the past few years. Choi et. al. find that
the zero-mode contributions in principle depend on the form of the vector-meson vertex
\( \Gamma^\mu = \gamma^\mu - (p_v - 2k)^\mu / D; \) the form factor \( f(q^2) \) is found to be free from a zero-mode contribution if the denominator \( D \) contains a term proportional to the light-front energy \( (k^-)^n \) with a power \( n > 0 \). Analogs of \( \Gamma^\mu \) occur for the spinor parts of the nucleon current matrix element, as we shall discuss below.

In this paper we calculate the electromagnetic form factors of the nucleon, taken as a scalar and axialvector diquark bound state, and investigate the zero-mode contributions in the quark-diquark model.

In the following section, the nucleon wave function in the diquark model and an effective quark-diquark-nucleon effective Lagrangian are given. Then the extraction of the nucleon form factors from the current matrix elements are given in Section 3. The discussion of the zero mode is given in Section 4. The numerical results for the form factors in the frame where the virtual photon momentum \( q^+ = 0 \) are given in Section 5. The summary and discussion are given in the last section.

2 Wave Function of the Nucleon

In the static case the nucleon wave function with \( SU(6) \) symmetry can be written as

\[
\frac{1}{\sqrt{2}} (\chi_\rho \Phi_\rho + \chi_\lambda \Phi_\lambda) \psi_0, \tag{1}
\]

where \( \chi_\rho \Phi_\rho \) corresponds to the scalar-isoscalar component of the nucleon and \( \chi_\lambda \Phi_\lambda \) to the vector-isovector quark pair part. This wave function can be rewritten in quark-diquark form as

\[
|p^+\rangle = \frac{1}{\sqrt{18}} [(2V_{11}^+ d^\uparrow - \sqrt{2}V_{11}^- d^\uparrow - \sqrt{2}V_{10}^+ u^\uparrow + V_{10}^- u^\uparrow) \sin \alpha + 3S_{00} u^\uparrow \cos \alpha], \tag{2}
\]

where the \( V \) are axial-vector diquark amplitudes and \( S \) is the scalar diquark amplitude. The explicit form of the \( V \) and \( S \) is given in Ref. [22]. If we set \( \alpha = \pi / 4 \), then Eq. (2) reduces to Eq. (1), that is, the permutation symmetry is restored.
Following KW [13], the two parts of the wave function in the nonstatic case become

\begin{align*}
\chi_\rho \Phi_\rho &\rightarrow N_S \left( \bar{q}_1 \gamma_5 \tau_2 C \bar{q}_2^T \right) (\bar{q}_3 N_\lambda) ; \\
\chi_\lambda \Phi_\lambda &\rightarrow N_V \left( \bar{q}_1 \gamma^\mu \tau \tau_2 C \bar{q}_2^T \right) \cdot (\bar{q}_3 \gamma_\mu \tau \gamma_5 N_\lambda), \tag{3}
\end{align*}

where a sum over quark permutations is implied. Here \( C = \gamma^2 \gamma^0 \) is the charge conjugation matrix, \( N_S, N_V \) are the normalization constants of the diquark wave functions. Also \( N_\lambda = [P]u_\lambda \) with \( u_\lambda(P) \) the Dirac-spinor of the nucleon and \( [P] \) denotes \( \gamma \cdot P + m_N \) with \( m_N \) the nucleon mass, \( q_i = [P]u_i \) with \( u_i \) the quark spinors; a factor \( [P] \) with each quark spinor appears with the Melosh rotation of the nonrelativistic spin wave function to the light cone. This is derived in detail in Sect. IV.A of Ref. [23].

Here we recognize \( \bar{q}_1 \gamma_5 \tau_2 C \bar{q}_2^T \) and \( \bar{q}_1 \gamma^\mu \tau \tau_2 C \bar{q}_2^T \) as the scalar and axialvector diquarks from the NJL model [24]. Thus, the nucleon wave function can be written as (see the initial quark-diquark vertex in Fig. 1)

\begin{align*}
|N\rangle &= f_S \varphi_S(x, k_{\perp}) \mathcal{O}_S u_\lambda(P) + f_V \varphi_V(x, k_{\perp}) \epsilon^i_\alpha \bar{\pi}(k) \mathcal{O}_V^\alpha \gamma_5 u_\lambda(P), \tag{4} \\
\mathcal{O}_S &= 1, \quad \mathcal{O}_V^\alpha = (\gamma^\alpha + P^\alpha/M)\gamma_5, \tag{5}
\end{align*}

where \( \varphi_S \), and \( \varphi_V \) are the scalar and axialvector radial diquark wave functions and \( \epsilon_\alpha \) is the axialvector diquark polarization vector, and \( k^\mu \) is the quark four-momentum. If coupling constants \( f_S \) and \( f_V \) is free completely, there will be a common factor, which will give a non-one charge of proton, so they must be constrained to give correct charge of proton. Our wave function is similar to that of Kroll et al. [17] and its basic diquark properties are consistent with those of other authors [24, 25].

From Eq. (4) we can introduce an effective Lagrangian for the quark-diquark-nucleon coupling as in Ref. [15] for the three-quark model

\begin{align*}
\mathcal{L}_{qDN} &= f_S \bar{\psi}(k_1) \mathcal{O}_S \psi_N(P) \phi_S(k_2) \Lambda_S + f_V \bar{\psi}(k_1) \mathcal{O}_V^\alpha(k_2) \frac{\tau^i}{\sqrt{3}} \psi_N(P) \cdot \phi_V^\alpha(k_2) \Lambda_V, \tag{6} \\
\mathcal{O}_S &= 1, \quad \mathcal{O}_V^\alpha = (\gamma^\alpha + i\partial^\alpha/M)\gamma_5, \tag{7}
\end{align*}
where $\psi$, $\psi_N$, $\phi_S$ and $\phi_V^\alpha$ are the quark, nucleon and scalar and axialvector diquark fields, respectively. The vertex functions $\Lambda_S$ and $\Lambda_V$ in Eq. (6) will be connected with the wave functions $\varphi_S(x, k_\perp)$ and $\varphi_V(x, k_\perp)$ later. In this work we adopt the same form for $\Lambda_S$ and $\Lambda_V$.

The Feynman rules of the diquark are presented in Table 1.

Table 1: Feynman rules of the diquark.

|       | $S$ | $V$ | $S\gamma S$ | $V\gamma V$ |
|-------|-----|-----|-------------|-------------|
| $S$   | $S\gamma S$ | $V\gamma V$ |
| $V$   | $V\gamma V$ |
| $S\gamma S$ | $\frac{i}{p^2-m_S^2} - \frac{i\delta_{ij}g_{\alpha\beta}}{p^2-m_V^2}$ | $\frac{i}{\sqrt{3}} \Gamma_{S\gamma S}^\mu$ | $Tr[\tau_i e_q \tau_j] \Gamma_{V\alpha\beta}^\mu$ |
| $V\gamma V$ | $i f_S$ | $f_V \mathcal{O}_V^\alpha \frac{\tau_i}{\sqrt{3}}$ | $rac{1}{3} \Gamma_{S\gamma S}^\mu$ |

The coupling of the virtual photon to a diquark is given as [17, 25]

$$\Gamma_S^\mu = -i(p_1 + p_2)\gamma^\mu, \quad (8)$$

$$\Gamma_{V\alpha\beta}^\mu = i \left( g_{\alpha\beta}(p_1 + p_2)^\mu - (1 + \kappa)g_{\alpha}^\mu(p_1 - p_2)\beta - (1 + \kappa)g_{\beta}^\mu(p_2 - p_1)\alpha \right) \quad (9)$$

with $\kappa$ the anomalous contribution of the magnetic moment of the axialvector diquark, $p_2$ the final momentum and $p_1$ the initial momentum of the diquark. Here, we have used the 't Hooft-Feynman gauge parameter $\xi = 1$, as in Keiner [25], which ensures the zero charge of the neutron.

### 3 Form Factors of the Nucleon

The electromagnetic form factors of the nucleon are conventionally defined in terms of the electromagnetic current matrix elements,

$$\langle N\lambda'|J^\mu|N\lambda \rangle = -ie \bar{u}_\lambda'(P') \left[ \gamma^\mu F_1(q^2) + \frac{1}{4m_N} q_\nu [\gamma^\nu, \gamma^\mu] F_2(q^2) \right] u_\lambda(P) \quad (10)$$

with $q = P' - P$, $P'$ and $P$ the final and initial total momentum.
Figure 1: Feynman diagrams for the nucleon current in case a photon strikes the quark. The case of a photon striking the diquark is analogous.

The form factors $F_i(q^2)$ can be determined from the $J^+$ matrix elements,

$$F_1 = \frac{m_N}{\sqrt{p_i^+ p_f^+}} J_{1\uparrow\uparrow}^+ + \frac{m_N}{\sqrt{p_i^+ p_f^+}} \frac{m_N q^{+2}_L}{Q_L^3} J_{1\downarrow\downarrow}^+,$$
$$F_2 = \frac{4 m_N^2 \sqrt{p_i^+ p_f^+}}{Q_L^3} J_{1\downarrow\uparrow}^+, \quad (11)$$

where $Q_L^3 = (p_L^+ p_f^+ + p_L^- p_f^-) q^+ - 2 p_f^+ p_f^- q_L$. If $q^+ = 0$, the form factor is obtained in the so-called Drell-Yan-West frame.

The Sachs form factors are

$$G_E(q^2) = F_1(q^2) + \frac{q^2}{4 m_N^2} F_2(q^2),$$
$$G_M(q^2) = F_1(q^2) + F_2(q^2). \quad (12)$$

The form factor problem is now reduced to calculating the matrix elements

$$J_{\lambda\lambda'}^+ = \int \frac{d^4 k}{(2\pi)^4} \frac{\Lambda(k_2') S_{\lambda\lambda'}^+ \Lambda(k_2)}{D(k_2') D(k) D(k_2)}, \quad (13)$$

as shown in Fig. 1. Here $\Lambda(k_2')$ and $\Lambda(k_2)$ are the vertex functions $\Lambda_S$ or $\Lambda_V$ given in Eq. 3, and

$$D(k) = k^2 - m_1^2 + i\varepsilon,$$
$$D(k_2') = k_2'^2 - m_2^2 + i\varepsilon = (P' - k)^2 - m_2^2 + i\varepsilon,$$
$$D(k_2) = k_2^2 - m_2^2 + i\varepsilon = (P - k)^2 - m_2^2 + i\varepsilon. \quad (14)$$
The spinor part $S_{\lambda \lambda'}^+ \gamma^\mu$ can be obtained from the Feynman rules in Table 1 and Fig. 1. If a photon strikes the quark, the quark is labeled as particle 2, with the spinor parts

$$S_{qS}^\mu = -i \langle f | [k_2'] e_q \gamma^\mu [k_2] | i \rangle,$$

(15)

$$S_{qV}^\mu = n \delta_{ij} g_{\alpha \beta} \langle f | \bar{O}_V^{\alpha \tau_i} \sqrt{3} [k_2'] e_q \gamma^\mu [k_2] O_V^{\beta \tau_j} \sqrt{3} | i \rangle,$$

(16)

in the numerator of Eq. (13) for the current matrix element, where we use the notation $[p] = \gamma \cdot p + m$, as earlier for $[P]$, and $\bar{O} = \gamma^0 O^\dagger \gamma^0$.

If the virtual photon strikes the diquark, the diquark is labeled as particle 2 with the spinor parts

$$S_{DS}^\mu = \langle f | [k] | i \rangle \frac{1}{3} \Gamma_S^\mu,$$

(17)

$$S_{DV}^\mu = \langle f | \bar{O}_V^{\alpha \tau_i} \sqrt{3} [k] O_V^{\beta \tau_j} \sqrt{3} | i \rangle Tr[\tau_i e_q \tau_j] \Gamma_V^{\alpha \beta},$$

(18)

with $\Gamma_S^\mu, \Gamma_V^{\alpha \beta}$ from Eqs. (8-9). Compared with the three-quark model, only Fig. 1(a) and 1(d) of Ref. [15] remain in the diquark model because we regard the diquark as a pointlike particle so that such details of its internal structure are no longer relevant.

From an explicit calculation we recognize the same isospin parts as in Kroll et al. [17]. For example, the isospin part of Eqs. (15-18) for the proton is $e_u, e_u + 2e_d, e_{ud}, e_{ud} + 2e_{uu}$ respectively, where $e_{ud}, e_{dd}, e_{uu}$ are the electric charges of the diquarks as given in Ref. [22].

4 Zero Mode Analysis

In order to search for zero modes, we now integrate the current matrix element over $k^-$ in Eq. (13) in the valence and nonvalence regions corresponding to Fig. (b) and (c) In frame where $q^+ > 0$. 

4.1 Integration in the Valence Region

In the valence region, characterized by $0 < k^+ < P^+$, there is only one pole $k^- = k^-_{on} = (m_1^2 + k_1^2 - i)/k^+$ located in the lower half of the complex $k^-$-plane. Thus, we can integrate over $k^-$ by Cauchy’s formula which yields

$$ J_{\lambda}^{\pm} = \int \frac{d^4k}{(2\pi)^4} \frac{h(k_2')S^+_{\lambda\lambda'}h(k_2)}{(k_2'^2 - m_2^2 + i\epsilon)(k_2'^2 - \Lambda^2 + i\epsilon)(k_2'^2 - m_2^2 + i\epsilon)(k_2'^2 - \Lambda^2 + i\epsilon)(k_2'^2 - m_2^2 + i\epsilon)}.$$  

where

$$M_0'^2 = \frac{k_1'^2 + m_1'^2}{x_2} + \frac{k_2'^2 + m_2'^2}{x_2}, \quad M_0'^2 = \frac{k_1'^2 + m_1'^2}{x_2} + \frac{k_2'^2 + m_2'^2}{x_2}, \quad M_0'^2 = \frac{k_1'^2 + \Lambda^2}{x_2} + \frac{k_2'^2 + \Lambda^2}{x_2}, \quad M_0'^2 = \frac{k_1'^2 + \Lambda^2}{x_2} + \frac{k_2'^2 + \Lambda^2}{x_2}$$  

with $x = k^+/P^+$, $x_2 = 1 - x$, $x' = k^+/P' = x/\alpha$, $x_2' = 1 - x'$, $k_\perp = k_\perp - x'q_\perp$.

Here we adopt $P_\perp = 0$, and $M_f$ and $M_i$ are the mass of the initial and final states, i.e. the nucleon mass. We use a Pauli-Villars regulator $\Lambda$ in $\Lambda(k_2) = h(k_2)/(k_2^2 - \Lambda^2 + i\epsilon)$, where $\Lambda$ plays a role of momentum cutoff, to make the integration in Eq. (19) and Eq. (24) below finite [8, 21]. From Eq. (19) we see that, if all the poles arising from $\Lambda(k_2)$, such as more Pauli-Villars fields, lie in the upper half of the complex $k^-$-plane, the explicit form of $\Lambda(k_2)$ has no effect on the result in the valence region provided the definition as Eq. (22) are chosen to include all factor from $\Lambda(k_2)$ to definition. A more explicit discussion of Pauli-Villars fields can be found in Ref. [26] in detail.

We can connect the integration in the valence region obtained above to the form of Ref. [13] directly by using the spin sum

$$2m[p] = \sum_{\lambda} u_{\lambda}(p)\bar{u}_{\lambda}(p).$$  

For example,

$$\frac{S^\mu_{ss}}{\sqrt{x_2'x_2}} = -i\langle f | [k_2']_e q^\mu [k_2]_e | i \rangle \propto \bar{u}_f u(k_2')\frac{\bar{u}(k_2')}{\sqrt{k_2'^+}} q^\mu \frac{u(k_2)}{\sqrt{k_2^+}} \bar{u}(k_2) u_i,$$  

obviously is a current density $\frac{\bar{u}(k_2')}{\sqrt{k_2'^+}} q^\mu \frac{u(k_2)}{\sqrt{k_2^+}}$ between two vertices $\bar{u}_f u(k_2')$ and $\bar{u}(k_2) u_i$ in
Eq. 4. Then we let

$$h(k_2) \frac{1}{\sqrt{2(2\pi)^3}} \frac{1}{(m_N^2 - M_0^2)x_2(m_N^2 - M_\Lambda^2)} x_2^2 (m_N^2 - M_\Lambda^2) = \varphi(x_2, k_{2\perp}),$$

(22)

where $h(k_2)/[x_2(m_N^2 - M_\Lambda^2)]$ corresponds to $\Lambda(k_2)$ before integration. $\varphi(x_2, k_{2\perp})$ is just the usual radial wave function in eq. (4). So if we use a $\Lambda(k_2)$ without pole, our definition of $\varphi(x_2, k_{2\perp})$ is the same as the one used by Araujo et. al. [15]. And only with such a definition do the $x, x_2, x_2'$ that are left give an invariant phase-space volume element, $d\Gamma = d^2k_\perp d^2k_{2\perp} \delta(k_\perp + k_{2\perp}) dx dx_2 / (x_2 x - 1)$ in the Drell-Yan-West frame [13].

We find for the valence region results similar to Ref. [13]. In the three-quark model the results with the integration as in this work are the same as the results in Ref. [13] because all terms with $[p]$ are on-shell so that they can be evaluated as $\sum \lambda u_\lambda(p)\bar{u}_\lambda(p)$.

In the $q^+ = 0$ frame, that is, the Drell-Yan-West frame, $J_{\lambda\lambda'}^+$ can be obtained by the limit

$$J_{\lambda\lambda'}^{+DYW} = \lim_{\alpha \to 1} J_{\lambda\lambda'}^+.$$

(23)

In the valence region corresponding to Fig. 1 (b), when $\alpha \to 1$, then $x' \to x, x_2' \to x_2$, and so on. We then find that $J_{\lambda\lambda'}^{+DYW}$ is obtained from the wave function in Eq. (2) and the free quark current density $\frac{\bar{u}(k_2')^\mu u(k_2)}{\sqrt{k_2^+} \sqrt{k_2'^+}}$ or the diquark current in Eqs. (8, 9) as usual. From Eq. (19) we see that, if a pole of $\Lambda(k_2)$ arises from $(k_2^2 - m^2 + i\epsilon)$, where $m = m_2, \Lambda, \text{ etc}$, it will lie in the upper half of the complex $k^-$-plane. Thus, the explicit form of $\Lambda(k_2)$ has no effect on the integration in the valence-quark region provided the diquark structure function $h(k_2)$ is chosen appropriately, as will be discussed below.

4.2 Integration in the Nonvalence Region

The nonvalence region is characterized by $P^+ < k_+ < P'^+$ kinematically. We can recognize Fig. 1 (c) as a pair creation process in the nonvalence region from the orientation of the momentum. Thus, the nonvalence contribution corresponds to a Z-diagram in the three-quark model if the virtual photon strikes a quark and corresponds to a double Z-diagram.
if the virtual photon strikes the diquark. Thus dynamically, the nonvalence region can no longer be described by a wave function. In the Drell-Yan-West frame there is no such contribution to the $J^+$ current component, so it must vanish when $q^+ \to 0$, that is, $\alpha \to 1$. Otherwise there would be a so-called zero mode contribution.

In the nonvalence region there are two poles located in the lower half of the complex $k^-$-plane when we integrate the current density over $k^-$ by Cauchy’s formula to yield

$$
J_{\lambda}^{+} = \int \frac{d^4 k}{(2\pi)^4} \frac{h(k') S_{\lambda\lambda}^{+} h(k_2)}{(k_2^2 - m_2^2 + i\epsilon)(k_2^2 - \Lambda_2^2 + i\epsilon)(k_2^2 - m_1^2 + i\epsilon)(k_2^2 - \Lambda_1^2 + i\epsilon)(k_2^2 - m_1^2 + i\epsilon)}
$$

$$
= 2\pi i \int_1^\alpha dx \frac{d k_1^2}{2(2\pi)^2 x_{\perp}^2 x_{\parallel}^2 (k_2^2 - m_2^2)} \left\{ \frac{h(k') S_{\lambda\lambda}^{+} h(k_2)}{(M_f^2 - M_0^2)(q^2 - M_{m2m2}^2)(q^2 - M_{m2\Lambda})} \right\}, \tag{24}
$$

where $x'' = x_2/(1-\alpha)$, $k_1'' = k_1 - x'' q_\perp$, $M_{m2m2} = k_1''^2 + m_2^2$, $M_{m2\Lambda} = k_1''^2 + \Lambda_2^2$, $M_{\Lambda m2} = k_1''^2 + \Lambda_1^2$.

If there are more Pauli-Villars fields used, that is, $\Lambda(k_2) = h(k_2)/\Pi_i(k_2^2 - \Lambda_i^2 + i\epsilon)$, we can use relation

$$
\frac{1}{(k_2^2 - \Lambda_i^2 + i\epsilon)(k_2^2 - \Lambda_j^2 + i\epsilon)} = \frac{1}{\Lambda_i^2 - \Lambda_j^2} \left( \frac{1}{k_2^2 - \Lambda_i^2 + i\epsilon} - \frac{1}{k_2^2 - \Lambda_j^2 + i\epsilon} \right) \tag{25}
$$

to split the integration to the sum of integration as Eq. (24).

Now we connect the integration in the nonvalence region contribution to the form of Ref. [13]. In the valence region we used the spin sum $2m[p] = \sum_\lambda u_\lambda(p) \bar{u}_\lambda(p)$. However, in the nonvalence region we have $x_i = x - 1$ for initial part, so we should define $k_2$ as $k - P$, and not $P - k$ as $k_2$, that is, all $k_2$ should be replaced by $-k_2$. For example,

$$
\frac{S_{\mu\nu}}{\sqrt{x_2 x_2}} = -i \langle f| \frac{[k_2']}{\sqrt{x_2'}} e_q \gamma^\mu \frac{[-k_2]}{\sqrt{x_2}} |\bar{e}\rangle \propto \bar{u}_i(u(k_2') \bar{u}(k_2') \gamma^\mu \frac{v(k_2)}{\sqrt{k_2^2}} \frac{v(k_2)}{\sqrt{k_2^2}} \bar{u}(k_2)) u_i \tag{26}
$$

using the spin sum

$$
2m[-p] = - \sum_\lambda v_\lambda(p) \bar{v}_\lambda(p). \tag{27}
$$
Obviously, this is a current density \( \bar{u}(k'_2) e_q \gamma^\mu \bar{v}(k_2) \) between two vertices \( \bar{u}f u(k'_2) \) and \( \bar{v}(k_2)u_i \). Hence the nonvalence region contribution is just the virtual pair creation process.

Because \( h(k_2) \) functions in both regions are from the same one in Eq. (13) we can still define \( h(k_2) \) as

\[
\frac{1}{\sqrt{2(2\pi)^3 (m_N^2 - M_0^2)x_2 (m_N^2 - M_2^2)}} = \varphi(x_2, k_{2\perp}).
\]

(28)

However in Eq. (24) only \( (M_2^f - M_0^2) \) and \( (M_2^f - M_N^2) \) occur for the final state. The vertex in the initial state is the so-called no-wave function vertex which, via Eq. (7), has a spin part different from the nucleon wave function but still a radial nucleon wave function. This vertex was also studied in Refs. [27, 28, 29]. In the work by Choi et. al. they use BS equation to connect the nonwave function vertex with the normal vertex with a kernel. In our work, we do not make explicit calculation of nonvalence contribution, so we use the results without using BS equation directly.

### 4.3 Recognition of Zero Modes

Now we proceed to look for a zero mode of the nucleon, that is, the spin-1/2 case. The zero mode contribution is defined as in Ref. [21]

\[
J_{\lambda \chi}^Z.M. = \lim_{\alpha \to 1} J_{\lambda \chi}^Z.NV.
\]

(29)

No singularities, hence zero modes, arise at \( x = 0 \). To isolate the zero mode contribution at \( x = 1 \) or \( x = \alpha \), we first make a transformation of the longitudinal integration variable,

\[
J_{\lambda \chi}^Z.M. \sim \lim_{\alpha \to 1} \int_1^\alpha dx \frac{x(\alpha - x)[x''(1 - x'')]^2}{xx''x''} h(k'_2)S_{\lambda \chi}^+ h(k_2)[...]
\]

\[
= \lim_{\alpha \to 1} \int_1^\alpha dx \frac{\alpha(\alpha - x)^2}{(1 - \alpha)^2} h(k'_2)S_{\lambda \chi}^+ h(k_2)[...]
\]

\[
= \lim_{\alpha \to 1} \int_0^1 dz \alpha(\alpha - 1)(1 - z)^2 h(k'_2)S_{\lambda \chi}^+ h(k_2)[...]
\]

(30)

where \( x = 1 - z(1 - \alpha) \) and \( x'' = (1 - x)/(1 - \alpha) = z \). So, both \( \alpha - x \) and \( 1 - x \) can give a factor \( 1 - \alpha \). Hereafter we identify \( \alpha - x \) and \( 1 - x \), when we consider the zero mode
contribution. Our analysis at the endpoint $1 - x$ will be based on counting powers of $1 - x$ in the integrand of Eq. (30).

Now we take

$$\varphi(x_2, k_{2\perp}) = \frac{1}{\sqrt{2(2\pi)^3}} \frac{\Lambda^2}{(m_N^2 - M_0^2)x_2(m_N^2 - M_0^2)}.$$ (31)

giving that $h(k_2) = \Lambda^2$ and $h(k'_2) = \Lambda^2$, which is just the form used by Choi et. al [21] as a starting point of their analysis, because this is the simplest choice and it means that there is no complicated internal structure in the vertex. Then the key to recognizing a zero-mode contribution is the spinor part $S^+_{\lambda\lambda'}$ in Eq. (13). Using Table IV in Ref. [13], it is easy to count $(1 - x)^n$ factors in $S^+_{\lambda\lambda'}$, with $(1 - x)$ coming from $x_2'$ or $x_2$. Every $[k_{2\text{on}}]$ or $[k'_{2\text{on}}]$ in $S^+_{\lambda\lambda'}$ leads to a term with the factor $(1 - x)^{-1}$, where $k_{2\text{on}}$ or $k'_{2\text{on}}$ are the momenta before and after the photon strikes the quark or diquark. Adjacent $[k_{2\text{on}}]$ and $[k'_{2\text{on}}]$ terms (that is, without another $[p_{\text{on}}]$ between them) give a single factor $(1 - x)$ because $A_{ik}^+ R$ or $K_{ik}^{R,L}$ in Eq. (A11) in Ref. [14] will give a $(1 - x)$ factor when $q \rightarrow 0$. The $\gamma^+$ and $\gamma^-$ between $[k_{2\text{on}}]$ and $[k'_{2\text{on}}]$ give a factor $(1 - x)$ and $(1 - x)^{-1}$, respectively, and $\gamma^+$ between $[k_{2\text{on}}]$ and $[k'_{2\text{on}}]$ gives a factor $(1 - x)$. In contrast, $\gamma^+$ and $\gamma^-$ in other cases and other $\gamma$’s, such as $\gamma_5$, $\gamma^L$, do not give rise to terms with factors $(1 - x)^n$.

Applying this power counting technique to the following example, when the spinor part is given by [21]

$$S^+_h = Tr[k'_2] \gamma^+(1 - \gamma_5)[k_2][k] \gamma_5 \epsilon^* \cdot \Gamma$$

$$= Tr[k'_{2\text{on}}] \gamma^+(1 - \gamma_5)[k_{2\text{on}}][k_{\text{on}}] \epsilon^* \cdot \Gamma \gamma_5$$

$$+ Tr[k'_{2\text{on}}] \gamma^+(1 - \gamma_5)[k_{2\text{on}}] \gamma^+ \epsilon^* \cdot \Gamma \gamma_5 \frac{1}{2} (k^- - k_{\text{on}}),$$ (32)

where $\Gamma^\mu = \gamma^\mu - (k + k')^\mu / D$ and $\epsilon_\alpha$ is a polarization vector, then in both lines after the second equality sign, the $(1 - x)^{-2}$ from $[k'_{2\text{on}}]$ and $[k_{2\text{on}}]$ is canceled by $\gamma^+(1 - \gamma_5)$ and the adjacent $[k'_{2\text{on}}]$ and $[k_{2\text{on}}]$ terms. For the $\epsilon^* \cdot \gamma$ part of $\epsilon^* \cdot \Gamma$ in the first line, there will be $(1 + x)$ and $(1 - x)^0$ from $\epsilon^* - \gamma^+ \gamma_5$ (or $\epsilon^* - \gamma^- \gamma_5$) and $\epsilon^* \cdot \gamma \gamma_5$ respectively. However $\epsilon^* \cdot (k + k')$
in the $D$ term will give a $(1 - x)^{-1}$ term. For the $\epsilon^* \cdot \gamma$ part of $\epsilon^* \cdot \Gamma$ in the second line, the
$\gamma^+$ from $\epsilon^* \cdot \Gamma$ disappears because of the other $\gamma^+$ using $\gamma^{+2} = \gamma^+$. The trace will let $[k'_{2n}]$
and $[k_{2m}]$ become adjacent again. So, the $(1 - x)^{-1}$ from $\frac{1}{2}(k^- - k_{2m})$ will be canceled. Then
there will be $(1 - x)$ and $(1 - x)^0$ from $\gamma^+ \epsilon^* \cdot \gamma \gamma_5$ and $\gamma^+ \epsilon^{\mu+} \gamma^- \gamma_5$, respectively. Thus, a zero
mode has to come from $D$ in $\Gamma^\mu$ in $(1 - x)^0$ order.

We now use the same method to analyze $S_{qS}, S_{qV}, S_{DS}, S_{DV}$ of Eqs. (35, 38).

For $S_{qS}$ the analog of Eq. (32) is $[k'_2]e_\gamma [k_2]$ from Eq. (15) in which the $(1 - x)^{-2}$ from $[k'_{2n}]$
and $[k_{2m}]$ is canceled by the $\gamma^+$ matrix element and the adjacent $[k'_{2n}]$ and $[k_{2m}]$ terms. The same applies to $S_{qV}$, where the analog of Eq. (32) is $O^\alpha_V [k'_2]e_\gamma [k_2]O^\beta_V$ from Eq. (16).

For $S_{DS}$ in Eq. (17), the instantaneous part of the term $[k]$, which comes from the quark propagator, is $\gamma^+(M^2 - M'^2)/2P^+$, which will give a $(1 - x)^{-1}$ factor. This $(1 - x)^{-1}$ factor will be canceled by that from $(k_2 + k'_2)^+$. Since there is already a factor $1 - \alpha$ in Eq. (30) there will be no zero mode from $S_{qS}, S_{qV}$ and $S_{DS}$.

Now we consider the spinor part $S_{DV}$ whose analog of Eq. (32) is $O^\alpha_V \frac{\tau_1}{\sqrt{3}} [k]O^\beta_V \frac{\tau_1}{\sqrt{3}} Tr[\tau_i e_\gamma \tau_j] \Gamma^\mu_{\alpha\beta}$
of Eq. (16) with $\Gamma^\mu_{\alpha\beta}$ from Eq. (9). As far as $g_{\alpha\beta}(p_1 + p_2)^\mu$ in $\Gamma^\mu_{\alpha\beta}$ is concerned, $(p_1 + p_2)^\mu$
cancel the $(1 - x)^{-1}$ factor from $[k]$ as in $S_{DS}$. However, for $(1 + \kappa)(g_{\alpha\beta} q_\beta - g_{\alpha\beta} q_\alpha)$, the factor
$(1 - x)^{-1}$ can not be canceled. Hence, if we remember that $h(k'_2)$ and $h(k_2)$ are constant, as
in Choi et. al. (21), there is a zero mode contribution in $(1 - x)^0$ order from $S_{DV}$.

From the above analysis, with a constant $h(k_2)$, there will be zero mode. Obviously, if we
choose $\varphi(x_2, k_{2\perp})$ to let $h(k_2)$ have a pole, that is, have a factor $(1 - x)^n$ with $n < 0$, there
will be a higher order zero mode while if $h(k_2)$ has a factor $(1 - x)^n$ with $n > 0$, which could
reduce $h(k_2)$ after the transformation in Eq. (30) to zero, the zero modes will disappear.

In the current matrix element, Eq. (31), the term $[k]$ comes from the quark propagator
whose instantaneous part is $\gamma^+(M^2 - M'^2)/2P^+$. Hence the corresponding instantaneous
spinor part

$$S_{DV_{inst}}^\mu = \langle f|O^\alpha_V \frac{\tau_1}{\sqrt{3}} [k]_{\text{inst}}O^\beta_V \frac{\tau_1}{\sqrt{3}} |i\rangle Tr[\tau_i e_\gamma \tau_j] \Gamma^\mu_{\alpha\beta}$$
\[
\tau^{it} \tau^{j} |N\rangle TR[\tau_i e_\tau \tau_j](M^2 - M'^2)/2P^{i+} \\
\cdot \bar{u}\chi\gamma_5(\gamma^\alpha + P^\alpha/M)\gamma^{+}(\gamma^\beta + P^\beta/M)\gamma_5 u_\lambda \\
\cdot \left\{ g_{\alpha\beta}(p_1 + p_2)^+ + (1 + \kappa)g_{\alpha}^+(k'_2 - k_2)_\beta - (1 + \kappa)g_{\beta}^+(k'_2 - k_2)_\alpha \right\} \\
\sim \bar{u}\chi[\gamma^+\gamma\cdot q - \gamma\cdot q\gamma^+ - \gamma^+q^2/M]u_\lambda \\
= C(1 + \kappa)(M^2 - M'^2) \begin{cases} 
P^+q^2/M^2, & \lambda' = 1/2, \lambda = 1/2 \\
2P^+q_L/M, & \lambda' = 1/2, \lambda = -1/2 
\end{cases}
\]

where \( C = \bar{u}\langle N | \tau^{it} \tau^{j} | N \rangle TR[\tau_i e_\tau \tau_j] \), \( M'^2 = M''^2 \) or \( M''^2 \). After \( \sim \), we have omitted all higher order terms of \((1 - \alpha)\) because they will vanish when \( q^+ \to 0 \). Hence there is clearly a zero mode in \( S_{DV} \), when \( q^+ \to 0 \).

### 4.4 Numerical Estimate of the Zero Mode

To give an estimate of the zero mode, in Fig. 2 we present the valence and nonvalence contribution of \( J_{DV}^+(Q^2)/J_{DV}^+(0) \) in the \( q^+ \to 0 \) limit with \( h(k'_2) \) and \( h(k_2) \) taken as constants using Eqs. (19, 24, 33). The valence contribution can be calculated using Eqs. (18, 19) directly. The nonvalence contribution can be obtained from Eqs. (18, 24) and the terms involving a zero mode have been given in Eq. (33). Here we use a mass \( m_q = 0.43 \text{ GeV} \) for the \( u \) (\( d \)) quark as in Refs. [19, 21] and a diquark mass of \( m_D = 0.6 \text{ GeV} \). We also adopt \( \Lambda = \Lambda' = 1.5 \text{ GeV} \). For both cases, the nonvalence contribution is visible. In the case \( \lambda = 1/2, \lambda' = 1/2 \) in Fig. 2(a), the nonvalence contribution increases with \( Q^2 \) and will cross the valence contribution at about \(-q^2 = 4 \text{ GeV}^2\). The non valence contributions will decrease slowly after about \( Q^2 = 8 \text{ GeV}^2 \). Hence the zero mode is far from negligible.

This power counting method we used is not required in the spin-1 meson case because the trace can be evaluated. However in the spin-1/2 case, this is difficult to do especially in the three quark model.
Figure 2: $J^+_{DV}(Q^2)$ in the valence and nonvalence regions divided by $J^+_{DV}(0)$ in the valence region. (a) is for the case $\lambda = 1/2, \lambda' = 1/2$ and (b) for the case $\lambda = 1/2, \lambda' = -1/2$. The full line is for the valence contribution, the dashed line for the nonvalence contribution.

5 Numerical Results

In Sec. 4, we have found that if the wave functions are chosen to let $h(k_2)$ and $h(k'_2)$ be constants, there will be a zero mode arising when $q^+ \to 0$, i.e. the Drell-Yan-West frame is used, which indicates that the nonvalence contribution must be considered. In the meson case, there are many works calculating a nonvalence contribution [27], [28], [29]. It would also be interesting to study the nonvalence contribution in the timelike region where $q^2 > 0$ so that we must use a frame with $q^+ > 0$. However in this paper, we focus on the spacelike region where $q^2 < 0$ so that Drell-Yan-West frame can be used. In case $h(k_2)$ and $h(k'_2)$ are not constant, from the analysis in Sec. 4 we can choose these wave functions to remove zero modes.

We calculate numerical results in the Drell-Yan-West frame defined by $q^+ = 0$. In this work we use for the wave functions $\varphi_S(x, k\perp)$ and $\varphi_V(x, k\perp)$ in Eq. 4 both the Gaussian parametrization

$$\varphi(x, k\perp) = N \exp(-M_0^2/8\beta^2),$$ (34)
and the negative power parametrization indicated by the form factor of Ref. [30].

\[ \varphi(x, k_{\perp}) = N (\omega^2 + M_0^2)^{-3.5}, \]  

(35)

where \( N \) is a normalization constant. Inserting the Eq. (33) to Eqs. (31, 24), from Eq. (30), we can find that the Gaussian form removes any zero mode because of \( M_0^2 \sim (x - 1)^{-1} \) (Here we note that in the nonvalence region \( x_{2i} = x - 1 \geq 0 \) and \( x_{2f} = \alpha - x \geq 0 \) as given in fig. 1(b), which ensures \( M_0^2 > 0 \) when \( q^+ \to 0 \)). The negative power shape wave functions will also remove the zero mode because \( h(k_2) \) and \( h(k_2') \) determined by this wave function will generate a \( (x - 1)^5 \) factor in Eq. (30).

The mass of the quark and diquark are chosen as \( mN/3 \) and \( 2mN/3 \), respectively, and \( \kappa = 1.6 \) as in Ref. [25] (in fact, the results are not sensitive to \( \kappa \)). From our calculations we find the results are more sensitive to the coupling constants \( f_S \) and \( f_V \). In the works using direct Melosh rotation [11, 31], the proton charge can be obtained automatically. However, as discussed in Sec 2, we let the scalar diquark coupling constant \( f_S \) determined by \( G_p(0) = 1 \). In Refs. [13, 32], they also normalized the wave function to charge one of proton. The free parameters left are \( f_V, \beta \) for the Gaussian type and \( \omega \) for negative power type wave functions. The zero charge of the neutron is guaranteed by the formula we used, that is, it is not fitted.

The parameters characterizing various fits are given in Table II.

Table 2: Parameters for Gaussian type wave function (Gau) and for negative power one (NP).

|       | \( m_q(MeV) \) | \( m_D(MeV) \) | \( \kappa \) | \( \beta(MeV) \) | \( \omega(MeV) \) |
|-------|----------------|----------------|-------------|----------------|-----------------|
| Gau   | 313            | 626            | 1.6         | 225            | ---             |
| NP    | 313            | 626            | 1.6         | ---            | 100             |

We present the charge and magnetic radii and magnetic moment of the nucleon with
different $f_V$ in Table III.

Table 3: Charge and magnetic radii and magnetic moments for different values $f_V$. The data are from Refs [33, 34, 35, 36].

| $f_V$(MeV) | $r_C^p$ (fm) | $r_M^p$ (fm) | $r_C^n$ (fm$^2$) | $r_M^n$ (fm) | $\mu_p(\mu_N)$ | $\mu_n(\mu_N)$ |
|------------|--------------|--------------|------------------|--------------|------------------|------------------|
| Gau I 80   | 0.804        | 0.834        | -0.262           | 0.861        | 2.225            | -1.009           |
| NP I 80   | 0.821        | 0.855        | -0.274           | 0.882        | 2.228            | -1.013           |
| Gau II 100 | 0.821        | 0.820        | -0.284           | 0.850        | 2.427            | -1.075           |
| NP II 100  | 0.836        | 0.837        | -0.295           | 0.868        | 2.428            | -1.078           |
| Gau III 130 | 0.854       | 0.796        | -0.325           | 0.832        | 2.814            | -1.201           |
| NP III 130 | 0.872       | 0.816        | -0.339           | 0.852        | 2.811            | -1.203           |
| Exp.       | --          | 0.870(8)     | 0.855(35)        | -0.116(2)    | 0.873(11)        | 2.793            | -1.913           |

As in the work of Kroll et al. [17], the results for the proton are better than those for the neutron. The charge and magnetic radii and magnetic moment are close to the data with $f_V = 130$ MeV. The magnetic moment of the neutron is too small, while the charge radius is too large, as is the case with other authors.

The electromagnetic form factors of the nucleon are shown in Fig. 3. Since the experimental magnetic form factors and the electric form factor of the proton are well described by a dipole fit

$$ G_E^p(Q^2) = G_M^p(Q^2)/\mu_p = G_M^n(Q^2)/\mu_n = F_D(Q^2) = \left(1 + \frac{Q^2}{m_D^2}\right)^{-2} \tag{36} $$

with $m_D^2 = 0.71$ GeV$^2$.

In Fig. 3 we present $G_E^p(Q^2)$, $G_M^p(Q^2)/\mu_p$ and $G_M^n(Q^2)/\mu_n$ divided by $F_D(Q^2)$ and the charge form factor of the neutron, $G_E^n(Q^2)/\mu_p$, with $f_V = 100$ MeV. The magnetic form factors of the proton and neutron are close to the dipole form with both wave functions while $G_M^n(Q^2)/\mu_n$ is lower than the dipole form, which is consistent with Jefferson Lab
data [37, 38, 39]. The magnetic form factor of the neutron is a bit too high.

In Fig. (4), we present the ratio \( \mu G_E/G_M \) with different \( f_V \) compared with recent data from Jefferson Lab. With all \( f_V \), the ratio will decrease in the higher \( Q^2 \) region and the falloff is too fast. There are similar results for the three-quark model in Ref [15]. The difference between the data and experiments may arise from the simplistic wave functions we use and omission of the intrinsic form factors of quarks and diquarks. In Fig. (4) we also find that the momentum transfer \( Q^2 \), where the electric form factor of the proton crosses zero, increases with increasing \( f_V \).

6 Summary and Discussion

We have given a relativistic diquark model of the nucleon using the KW method. We calculate the electromagnetic form factors of the nucleon taken as a scalar and axialvector diquark bound state. In this work we do not consider the effect of scalar-vector transitions and pion or gluon exchange effects, which may improve the numerical results. The charge and magnetic
Figure 4: The ratio $\mu G_E^p/G_M^p$. The thick and thin lines are for the Gaussian type and negative power type wave functions with $f_V = 80 \text{ MeV}$ (full), $100 \text{ MeV}$ (dashed), $130 \text{ MeV}$ (dotted) respectively. Experimental data are from Refs. [37, 38, 39].

radii and magnetic moment of the proton can be reproduced, while the magnetic moment of the neutron is too small. The shape of the form factors $G_M^p(Q^2)/\mu_p$ and $G_M^n(Q^2)/\mu_n$ can be reproduced also. The zero charge form factor of neutron at $Q^2 = 0$ (neutron charge) is guaranteed by the formalism we used.

There will be a zero mode in electromagnetic form factors of the nucleon arising from the instantaneous part of the quark propagator if we adopt $\varphi(x_2, k_{2\perp})$ to let $h(k_2)$ have $(1 - x)^n$ with $n \leq 0$. In our numerical calculation, we find that the zero mode cannot be neglected. Hence the quark-diquark-nucleon vertex functions must be chosen to remove the zero mode to ensure the correct results in the Drell-Yan-West frame.

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