Nucleon and pion electromagnetic form factors in a light-front constituent quark model

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Abstract

Nucleon and pion electromagnetic form factors are evaluated in the spacelike region within a light-front constituent quark model, where eigenfunctions of a mass operator, reproducing a large set of hadron energy levels, are adopted and quark form factors are considered in the one-body current. The hadron form factors are sharply affected by the high momentum tail generated in the wave function by the one-gluon-exchange interaction. Useful information on the electromagnetic structure of light constituent quarks can be obtained from the comparison with nucleon and pion experimental data.

The measurement of the electromagnetic (e.m.) form factors of hadrons represents a valuable tool for investigating in detail their internal structure. This fact has motivated a great deal of experimental and theoretical work, that will increase with the advent of new accelerator facilities, e.g. CEBAF, yielding unique information on the transition region from the non perturbative to the perturbative regime of QCD [1, 2]. Though the fundamental theory of the strong interaction, QCD, should be applied for describing hadron structure, the practical difficulties to be faced in the nonperturbative regime have motivated the development of effective theories, e.g. constituent quark (CQ) models, that in turn could provide useful hints to model approximations to the "true" field theory [3]. Aim of this letter is to apply our approach [4, 5], based on a relativistic CQ model, to the evaluation of the nucleon e.m. form factors in the spacelike region, keeping safe the good description already obtained for the pion form factor. Our model represents an extension of the one proposed in Refs. [6, 7], where a relativistic treatement of light CQ's is achieved by adopting the light-front formalism [8] and gaussian wave functions are assumed for describing the pointlike CQ's inside the nucleon (see also [9]). In particular we have considered: i) hadron wave functions which are eigenvectors of a light-front mass operator, constructed from the effective $q\bar{q}$ and qq-interaction of Refs. [10, 11], that reproduces a huge amount of energy levels; ii) the configuration mixing, due to the one-gluon-exchange (OGE) part of the effective interaction, leading to high momentum components and $SU(6)$ breaking terms in the hadron wave function; iii) Dirac and Pauli form factors for the CQ's, as suggested by their quasi-particle nature (cf. [12]), summarizing the underlying degrees of freedom. The comparison of our calculations with the experimental data on nucleon and pion form factors will phenomenologically constrain the e.m. structure of the light CQ's.

As known (cf. [8]), the light-front wave functions of hadrons are eigenvectors of a mass operator, e.g. $\mathcal{M} = M_0 + \mathcal{V}$, and of the non-interacting angular momentum operators j^2 and j_n, where the vector $\hat{n} = (0, 0, 1)$ defines the spin quantization axis. The operator M_0 is the free mass and the interaction term $\mathcal{V}$ is a Poincaré invariant. In this letter we briefly present the formalism for the nucleon only, since the relevant formulae for the pion can be found in [4]. For the three-quark system $M_0^2 = \sum_{i=1}^{3} \frac{q_{i\perp}^2 + m_i^2}{\xi_i}$, where m_i is the CQ mass, $\xi_i = p_i^+/P^+$ and $\vec{q}_{i\perp} = \vec{p}_{i\perp} - \xi_i \vec{P}_\perp$ are the intrinsic light-front variables. The subscript $\perp$ indicates the projection perpendicular to the spin quantization axis and the *plus* component of a 4-vector $p \equiv (p^0, \vec{p})$ is given by $p^+ = p^0 + \hat{n} \cdot \vec{p}$; finally $\tilde{P} \equiv (P^+, \vec{P}_\perp) = \tilde{p}_1 + \tilde{p}_2 + \tilde{p}_3$ is the light-front nucleon momentum and $\tilde{p}_i$ the quark one. In terms of the longitudinal momentum q_{in}, related to the variable ξ_i by $q_{in} = \frac{1}{2}\left(\xi_i - \frac{q_\perp^2 + m_i^2}{\xi_i M_0}\right)$, the free mass operator acquires a familiar form, viz.

$$M_0 = \sum_{i=1}^{3} \sqrt{m_i^2 + q_i^2} = \sum_{i=1}^{3} E_i \qquad (1)$$

with $\vec{q}_i \equiv (\vec{q}_{i\perp}, q_{in})$. Disregarding the color degree of freedom, the light-front nucleon wave

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function can be written as follows (see also [6, 7])

$$\langle\{\xi_i \vec{q_{i\perp}}; \nu_i \tau_i\}|\Psi_N^{\nu_N}\rangle \;=\; \sqrt{\frac{E_1 E_2 E_3}{M_0 \xi_1 \xi_2 \xi_3}} \sum_{\{\nu_i'\}} \langle\{\nu_i\}|\mathcal{R}^\dagger|\{\nu_i'\}\rangle\langle\{\vec{q_i}; \nu_i' \tau_i\}|\psi_N^{\nu_N}\rangle \tag{2}$$

where the curly braces { } mean a list of indices corresponding to $i = 1, 2, 3$; $\nu_i'(\tau_i)$ is the third component of the quark spin (isospin); $\mathcal{R}^\dagger = \left[\prod_{j=1}^{3} R_M^\dagger(\vec{q}_{j\perp}, \xi_j, m_j)\right]$ with $R_M(\vec{q}_{j\perp}, \xi_j, m_j)$ being the usual generalized Melosh rotation [13]. Disregarding the P and D waves (see below), the equal-time nucleon wave function $\langle\{\vec{q_i}; \nu_i' \tau_i\}|\psi_N^{\nu_N}\rangle$ is given by

$$\langle\{\vec{q_i}; \nu_i' \tau_i\}|\psi_N^{\nu_N}\rangle \;=\; \frac{1}{\sqrt{2}}\left[w_S(\vec{p}, \vec{k})\left(\Phi_{\nu_N \tau_N}^{00}(\{\nu_i' \tau_i\}) + \Phi_{\nu_N \tau_N}^{11}(\{\nu_i' \tau_i\})\right) \;+\;\right.$$

$$\left. w_{S'}^{s}(\vec{p}, \vec{k})\left(\Phi_{\nu_N \tau_N}^{00}(\{\nu_i' \tau_i\}) - \Phi_{\nu_N \tau_N}^{11}(\{\nu_i' \tau_i\})\right) \;+\; w_{S'}^{a}(\vec{p}, \vec{k})\left(\Phi_{\nu_N \tau_N}^{01}(\{\nu_i' \tau_i\}) + \Phi_{\nu_N \tau_N}^{10}(\{\nu_i' \tau_i\})\right)\right] \tag{3}$$

where $\vec{p} = \vec{q_1}$ and $\vec{k} = (\vec{q_2} - \vec{q_3})/2$ are the Jacobi coordinates for the three-quark system (cf. [7]) and the functions w_S and $w_{S'}^{s(a)}$ correspond to the S and S' (mixed-symmetry) waves, respectively. Finally, the spin-isospin function $\Phi_{\nu_N \tau_N}^{ST}(\{\nu_i' \tau_i\})$ is defined as follows

$$\Phi_{\nu_N \tau_N}^{ST}(\{\nu_i' \tau_i\}) = \sum_{M_S}\langle\tfrac{1}{2}\nu_1' S M_S|\tfrac{1}{2}\nu_N\rangle\,\langle\tfrac{1}{2}\nu_2'\tfrac{1}{2}\nu_3'|S M_S\rangle \sum_{M_T}\langle\tfrac{1}{2}\tau_1 T M_T|\tfrac{1}{2}\tau_N\rangle\,\langle\tfrac{1}{2}\tau_2\tfrac{1}{2}\tau_3|T M_T\rangle \tag{4}$$

The normalization of (3) is $\sum_{\{\nu_i \tau_i\}}\int\prod_{j=1}^{3}d\vec{q_j}\,\delta(\vec{q_1} + \vec{q_2} + \vec{q_3})\,|\langle\{\vec{q_i}, \nu_i \tau_i\}|\psi_N^{\nu_N}\rangle|^2 \;=\; 1$. The equal-time state $|\psi_N^{\nu_N}\rangle$ is an eigenvector of the transformed mass operator $M = \mathcal{R}M\mathcal{R}^\dagger = \mathcal{R}M_0\mathcal{R}^\dagger + \mathcal{R}V\mathcal{R}^\dagger$. Since the free-mass commutes with the Melosh rotation one has $M = M_0 + V$, where the interaction $V = \mathcal{R}V\mathcal{R}^\dagger$ has to fulfil the proper invariance properties, namely: i) no dependence upon the total momentum and the centre of mass coordinates, and ii) invariance upon rotations (cf. [8]). We have chosen M equal to the effective hamiltonian proposed by Capstick and Isgur (CI) [11]. Thus, our equal-time baryon states $|\psi_{qqq}\rangle = \sum_{\{\nu_i \tau_i\}}(\prod_{j=1}^{3}|\tfrac{1}{2}\nu_j\rangle|\tfrac{1}{2}\tau_j\rangle)\langle\{\vec{q_i}; \nu_i \tau_i\}|\psi_{qqq}\rangle$ are eigenvectors of the following mass operator

$$H_{qqq}\,|\psi_{qqq}\rangle \;\equiv\; \left[\sum_{i=1}^{3}\sqrt{m_i^2 + q_i^2} + \sum_{i\neq j=1}^{3} V_{ij}\right]|\psi_{qqq}\rangle \;=\; M_{qqq}\,|\psi_{qqq}\rangle \tag{5}$$

where M_{qqq} is the mass of the baryon, and V_{ij} the CI effective qq potential, which is composed by a OGE term (dominant at short separations) and a linear confining term (dominant at large separations). It should be pointed out that the relativistic mass operator (5) reproduces a large set of hadron energy levels [10, 11] and generates a huge amount of configuration mixing, due to the presence of the OGE part of the interaction. We have evaluated the nucleon form factors using in Eq.(2) the eigenstate of Eq.(5), whereas in the current literature only the effects of the confinement scale have been considered, through gaussian or power-law wave functions [6, 7, 9]. As in Refs. [10, 11], the values $m_u = m_d = 0.220\ GeV$ have been adopted.

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The wave equation (5) has been solved by expanding the wave functions $w_S(\vec{p}, \vec{k})$ and $w_{S'}^{s(a)}(\vec{p}, \vec{k})$ onto a (truncated) set of harmonic oscillator (HO) basis states (details will be given elsewhere [14]) and applying to the Hamiltonian H_{qqq} the Rayleigh-Ritz variational principle. We have checked that the convergence for all the quantities considered in this work can be reached including in the expansion all the basis states up to 20 HO quanta. In particular, the calculated nucleon mass is $M_N = .940\ GeV$ and the percentages of the various waves are: $p_S = 98.1$, $p_{S'} = 1.7$ and $p_D = 0.2$. The P wave has been neglected since it does not couple to the main component of the wave function. As in the case of pseudoscalar [4] and vector [5] mesons, the OGE part of the qq-interaction, determining the hyperfine splitting of hadron mass spectra, generates high momentum components in the baryon wave function [14, 15].

In what follows we will present the calculations of the nucleon and pion e.m. form factors performed by using the eigenvectors (2) and the one-body component of e.m. current, which for the nucleon is given by

$$\mathcal{I}_N^\nu = \sum_{j=1}^{3} I_j^\nu = \sum_{j=1}^{3} \left(e_j \gamma^\nu f_1^j(Q^2) \ + \ i\kappa_j \frac{\sigma^\nu_{\ \mu} q^\mu}{2m_j} f_2^j(Q^2) \right) \tag{6}$$

where $\sigma^\nu_{\ \mu} = \frac{i}{2}[\gamma^\nu, \gamma_\mu]$, e_j is the charge of the j-th quark, κ_j the corresponding anomalous magnetic moment, $f_{1(2)}^j$ its Dirac (Pauli) form factor, and $Q^2 \equiv -q \cdot q$ the squared four-momentum transfer. In the light-front formalism, see e.g. Ref. [6], the spacelike e.m. form factors are related to the plus component of the e.m. current evaluated in a frame where $q^+ = 0$; such a choice allows to suppress the contribution of the pair creation from the vacuum [1, 16]. For the nucleon one has

$$\langle \Psi_N^{\nu'_N} | \mathcal{I}_N^+ | \Psi_N^{\nu_N} \rangle = \mathcal{F}_1^N(Q^2)\delta_{\nu'_N \nu_N} - i\langle \tfrac{1}{2}\nu'_N | \sigma_2 | \tfrac{1}{2}\nu_N \rangle \sqrt{\eta_N} \mathcal{F}_2^N(Q^2) \tag{7}$$

where σ_2 is a Pauli matrix, $\mathcal{F}_{1(2)}^N$ the Dirac (Pauli) form factor of the nucleon and $\eta_N = Q^2/(4M_N^2)$. The comparison with the experimental data will be presented in terms of the Sachs form factors, i.e. $G_E^N = \mathcal{F}_1^N - \eta_N \mathcal{F}_2^N$ and $G_M^N = \mathcal{F}_1^N + \mathcal{F}_2^N$. The explicit expression for the nucleon form factors, including the contributions from S and S' waves, and for the pion one, including the contribution from the Pauli quark form factor not considered in [4], will be given elsewhere [15]. As a check, we have repeated the calculations of the nucleon form factors of Refs. [6, 7], where a simple S-wave gaussian function was adopted. It should be pointed out that the numerical calculations involve multifold integrations, carried out through a Monte Carlo routine [17].

In order to investigate the sensitivity of the form factor upon the high momentum tail of the nucleon wave function, we have calculated the nucleon form factors assuming pointlike quarks (i.e. $f_1^i = 1$ and $\kappa_i = 0$) and using different nucleon wave functions $\psi_N^{(CI)}$, $\psi_N^{(si)}$ and $\psi_N^{(conf)}$. These wave functions are obtained as solutions of Eq. (5) by considering the full interaction $V_{(CI)}$, the spin-independent part of $V_{(CI)}$ and only its linear confining part, respectively. The results for the ratio G_E^p/G_D (with $G_D = 1/(1 + Q^2/0.71)^2$) are shown

in Fig.1. The following comments are in order: i) the configuration mixing generated by the OGE part of the interaction sharply affects the nucleon form factor even at low values of Q^2, as in the case of the pion [4], ii) the results of the calculations performed with $\psi_N^{(conf)}$ are similar to the ones obtained with a gaussian wave function, e.g., in Ref. [7]. It turns out that analogous results hold as well for the other e.m. form factors. In particular, if $\psi_N^{(conf)}$ and $\psi_N^{(CI)}$ are used and pointlike CQ's are assumed, the values of the nucleon magnetic moments result to be $\mu_{p[n]}^{(conf)} = 2.74\,[-1.60]$ and $\mu_{p[n]}^{(CI)} = 2.28\,[-1.19]$, respectively; this means that the configuration mixing leads to a sizeable underestimation of the nucleon magnetic moments ($\mu_{p[n]}^{exp} = 2.793\,[-1.913]$).

Once the high momentum components, generated by the full CI qq-interaction, are taken into account, the nucleon form factors sharply differ from the dipole prediction (cf. the solid line in Fig.1). A possible way to solve such a discrepancy is to assume a non trivial e.m. structure for extended CQ's (see e.g. [12, 18]). Thus, a simple parametrization of the isoscalar and isovector parts of the CQ form factors $f_{1(2)}^i$ has been introduced

$$f_1^{S(V)}(Q^2) = \frac{e_u f_1^u \pm e_d f_1^d}{(e_u \pm e_d)} = \frac{A_1^{S(V)}}{1 + Q^2\, B_1^{S(V)}} + \frac{1 - A_1^{S(V)}}{(1 + Q^2\, C_1^{S(V)}/2)^2}$$

$$f_2^{S(V)}(Q^2) = \frac{\kappa_u f_2^u \pm \kappa_d f_2^d}{(\kappa_u \pm \kappa_d)} = \frac{A_2^{S(V)}}{(1 + Q^2\, B_2^{S(V)}/2)^2} + \frac{1 - A_2^{S(V)}}{(1 + Q^2\, C_2^{S(V)}/3)^3} \tag{8}$$

Within one-body approximation of the e.m. current the values of κ_u and κ_d can be fixed by the request of reproducing the experimental nucleon magnetic moments; in particular, using $\psi_N^{(CI)}$ we have obtained $\kappa_u = 0.085$ and $\kappa_d = -0.153$. Differently, the 12 constants $A_{1(2)}^{S(V)}$, $B_{1(2)}^{S(V)}$ and $C_{1(2)}^{S(V)}$ have been estimated through a standard minimization procedure, where the experimental form factors of both nucleon and pion in a wide range of momentum transfer, reaching for the proton $Q^2 \approx 30(GeV/c)^2$, have been considered. This wide range can be justified by the phenomenological attitude adopted in this letter and is also inspired by the fact that the onset of the perturbative QCD is not definitively localized [1, 2].

Our calculations (solid lines) are compared with the nucleon and pion data in Figs. 2 and 3, respectively. It should be pointed out that our results have been obtained in the framework of the one-body approximation for the hadron e.m. current, namely by disregarding the two-body currents necessary for fulfilling both the gauge and rotational invariances (see Ref. [8]). However, Figs. 2 and 3 clearly show that an overall agreement with the data can be achieved by assuming an effective one-body current that contains CQ's with a structure. The corresponding CQ form factors, $f_{1(2)}^u$ and $f_{1(2)}^d$, are presented in Fig. 4 and exhibit a difference between the u- and d-quark e.m. structure. In order to illustrate the role played by such a difference, we have also presented the results (dotted lines in Fig. 2) obtained by assuming $f_{1(2)}^S = f_{1(2)}^V$, with $f_{1(2)}^V$ given by the previous procedure, and the values of the anomalous magnetic moments fixed as follows: $\tilde{\kappa}_u/e_u = \tilde{\kappa}_d/e_d = \kappa_V = \kappa_u - \kappa_d = 0.238$. This prescription, which assumes the same e.m. structure for u and d quarks, does not change the prediction for the pion and helps to illustrate the relevance of the differences between the u and d quarks for explaining the nucleon data with accuracy. The nucleon charge form

factors, corresponding to the full calculations (solid lines in Fig. 2), have the following slopes at $Q^2 = 0$: $dG_E^{p(th)}(Q^2)/dQ^2 = -2.83 \pm .20 (c/GeV)^2$ $[dG_E^{p(exp)}(Q^2)/dQ^2 = -3.0 \pm .09 (c/GeV)^2$ [19]] and $dG_E^{n(th)}(Q^2)/dQ^2 = 0.33 \pm .03 (c/GeV)^2$ $[dG_E^{n(exp)}(Q^2)/dQ^2 = 0.50 \pm .015 (c/GeV)^2$ [19]], respectively. The theoretical error bars are estimated from the statistical uncertainties of the Monte Carlo numerical integration procedure.

As to the pion, due to the constraints imposed by the nucleon data, in particular by G_E^n, the overall quality of the fit is a little bit lower ($r_{ch}^{\pi(th)} = 0.71 fm$ while $r_{ch}^{\pi(exp)} = 0.660 \pm 0.024\ fm$ [29]) than the one obtained in Ref. [4], where only the pion was considered. It should be pointed out that, if the data for G_E^n are disregarded in the minimization procedure, an impressive agreement for all the remaining form factors is obtained, but the predicted G_E^n becomes quite small [15].

Finally, it is worth noting that the CQ form factors shown in Fig. 4 yield a quark radius, defined as $< r_1^{u(d)} >^2 = -6\ df_1^{u(d)}(Q^2)/dQ^2$ at $Q^2 = 0$, equals to $< r_1^{u(d)} >= 0.51 fm\ \ (0.42 fm)$; such values are similar to the ones obtained in Refs. [12, 32] and from our exploratory analysis of the pion data [4] (cf. also dotted line in Fig. 4).

In conclusion, the picture stemming from our analysis points to a description of the hadron form factors, in a wide range of momentum transfer, in terms of effective quarks, having a non trivial e.m. structure. For the calculation of the pion and nucleon form factors we have adopted the eigenfunctions of a light-front mass operator, reproducing a large set of hadron energy levels (see Eq.(5)), and a one-body approximation for the e.m. current, containing CQ form factors. Within this framework, the existing pion and nucleon data phenomenologically constrain the CQ form factors and this fact will allow parameter-free calculations of other hadron form factors in the u-d sector. The application of our approach to the calculation of the magnetic form factor of the $N - \Delta$ transition [15] as well as the evaluation of the corresponding angular condition [8], yielding an estimate of the effects of two-body e.m. currents, are in progress.

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Figure Captions

Fig. 1. The proton form factor G^p_E/G_D vs. Q^2 for different choices of the proton wave function, and assuming a pointlike structure for the CQ (i.e. $f_1^{u(d)} = 1$ and $\kappa_{u(d)} = 0$). Solid line: G^p_E/G_D calculated through $\psi_N^{(CI)}$, eigenfunction of Eq.(5) corresponding to the full interaction $V_{(CI)}$ of Ref. [11]; dashed line: G^p_E/G_D calculated with $\psi_N^{(si)}$, corresponding to the spin-independent part of the full interaction $V_{(CI)}$; dotted line: G^p_E/G_D calculated by using $\psi_N^{(conf)}$, corresponding to the linear confining part of the full interaction $V_{(CI)}$. For comparison the result of Ref. [7] has also been shown (dot-dashed line).

Fig. 2. - a) The proton form factor G^p_E/G_D vs. Q^2. Solid line: G^p_E/G_D obtained by using: i) the wave function $\psi_N^{(CI)}$ corresponding to the full interaction $V_{(CI)}$ of Ref. [11], ii) the nucleon e.m. current of Eq. (6), and iii) the CQ form factors of Eq.(8) (see text). Dotted line: the theoretical G^p_E/G_D calculated assuming $f_{1(2)}^S = f_{1(2)}^V$ and $\tilde{\kappa}_u/e_u = \tilde{\kappa}_d/e_d = \kappa_V = \kappa_u - \kappa_d = 0.238$ (see text). Experimental data are from Refs.[19] (full dots), [20] (open squares), [21] (open diamonds) and [22] (full squares). - b) The same of Fig. 2a, but for $G^p_M/(\mu_p\, G_D)$. Experimental data are from Refs.[19] (full dots), [20] (open squares), [21] (open diamonds) and [22] (full squares).- c) The same of Fig. 2 a, but for G^n_E. Experimental data are from Ref.[23] (full dots), corresponding to the analysis in terms of the N-N Reid soft core interaction, and from Ref. [24] (open squares). - d) The same of Fig. 2 a, but for $G^n_M/(\mu_n\, G_D)$. Experimental data are from Refs.[25] (open diamonds), [26] (full diamonds), [27] (full squares), [24] (full dots) and [28] (open squares).

Fig. 3. The charge form factor of the pion $F_\pi(Q^2)/F_{Mon}(Q^2)$ vs. Q^2, with $F_{Mon}(Q^2) = 1/(1 + Q^2/0.54)$. The theoretical curve has been obtained by using the wave function of the pion corresponding to the full interaction of Ref. [10], and the CQ form factors of Eq.(8). Experimental data are from Refs.[29] (open dots), [30] (open diamonds), and [31] (full squares).

Fig. 4. The constituent quark form factors, extracted from the analysis of the nucleon and pion form factors (see Figs. 2 and 3), vs. Q^2. Solid lines represent $f_1^u(Q^2)$ (thick) and $f_2^u(Q^2)$ (thin), respectively; dashed lines represent $f_1^d(Q^2)$ (thick) and $f_2^d(Q^2)$ (thin), respectively. For comparison $f_1^u(Q^2) = f_1^d(Q^2)$ obtained in Ref. [4] is also shown (dotted line).