The relativistic constituent quark model of low-energy quantum chromodynamics is found to yield a consistent picture of the electroweak structure of the nucleons. Notably, the electromagnetic and axial form factors of both the proton and the neutron can be described in close agreement with existing experimental data in the domain of low to moderate momentum transfers. For the theory it is mandatory to respect Poincaré invariance and to fulfill additional conditions like charge normalization. Here we present covariant predictions of the one-gluon-exchange and Goldstone-boson-exchange constituent quark models for the electroweak form factors of the nucleons and give a critical discussion of the results in view of the point-form spectator model employed for the electromagnetic and axial current operators.

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The explanation of the electroweak structure of the nucleons and of other baryon ground states still represents a formidable problem. Even though the theoretical framework appears to be well founded in the standard model of strong and electroweak interactions and a wealth of experimental data has been accumulated up till now, one has not yet reached a conclusive understanding of the electroweak and axial form factors, especially of the nucleons, at low and intermediate energies. Of course, the essential difficulties reside in the solution of quantum chromodynamics (QCD) outside the perturbative regime. Recently, calculations of the nucleon electroweak form factors have become available from relativistic constituent quark models (RCQMs). The pertinent covariant results have turned out remarkable in several respects. The RCQM assumes the nucleons to consist of three constituent quarks and describes their mass spectra in a Poincaré-invariant manner. The theory is formulated along relativistic Hamiltonian dynamics. Thus one works a-priori with a finite number of degrees of freedom rather than facing problems, such as regularization and truncation, in a field-theoretical approach. The symmetries of Lorentz invariance are strictly included by fulfilling the constraints of the Poincaré algebra. Similarly the essential properties of (nonperturbative) QCD, like the consequences of spontaneous breaking of chiral symmetry at low energies, can be implemented in the Hamiltonian or equivalently in the invariant mass operator defining the RCQM.

Elastic electroweak nucleon form factors have recently been studied in the framework of relativistic quantum mechanics. Specifically, working within the point form allows to calculate all the desired observables in a covariant manner. One has first produced the predictions of the Goldstone-boson-exchange (GBE) RCQM for the electroweak structure of the nucleons using spectator-model currents. The electromagnetic and axial nucleon form factors were readily found to be remarkably consistent with all experimental data at low momentum transfers. Here, we add the corresponding predictions of the one-gluon-exchange (OGE) RCQM by means of the relativistic version of the Bhaduri-Cohler-Nogami (BCN) model as parametrized in ref. The results are collected in Figures There a comparison is given to the predictions of the GBE RCQM, to the nonrelativistic impulse approximation (NRIA) and to the existing experimental data. One observes only minor differences between the RCQMs with different dynamics for the hyperfine interaction (the solid and dashed curves in Figures are practically indistinguishable, except for the neutron electric form factor ). The reason is that the relevant components of the nucleon ground-state wave functions are rather similar for both the OGE and GBE RCQMs. The differences between the two types of RCQMs become striking only for the excited states. From the form factor results in Figures it is also immediately evident that relativistic effects are of paramount importance in all respects. The NRIA fails completely.

In view of the existing results, one must ask why such a consistent picture of the electroweak structure of the nucleons can really come about by employing RCQMs; even more so since a simplified model of the electromagnetic and axial currents has been employed. Up till now the full many-body character of the electroweak currents cannot be tackled in a fully relativistic calculation. Rather one has been resorting to simplifications. The results exemplified here have been obtained within the so-called point-form spectator model (PFSM) for the electromagnetic and axial currents. The PFSM is characterized by the fact that the intermediate boson couples only to one of the constituent quarks in the baryon, while the momentum is transferred to the baryon as a whole. Among the available forms of relativistic quantum mechanics, the
point form is specific in the respect that the approximated current operators preserve their spectator character in all reference frames [11]. One uses a covariant (spectator-model) current operator and in addition further symmetry requirements, like translational and time-reversal invariance as well as charge normalization, are implemented in the construction.

The concrete expressions for the reduced matrix elements of the PFSM current between three-body states of quarks with individual momenta $p_i$ and spin projections $\sigma_i$ read

$$\langle p_1', p_2, p_3'; \sigma_1', \sigma_2', \sigma_3' | \hat{J}^{\mu}_{\text{spec}} | p_1, p_2, p_3; \sigma_1, \sigma_2, \sigma_3 \rangle = 3N \langle p_1', \sigma_1' | \hat{J}^{\mu}_{\text{spec}} | p_1, \sigma_1 \rangle \left(2p_{20} \delta^3(\vec{p}_2 - \vec{p}_2') + 2p_{30} \delta^3(\vec{p}_3 - \vec{p}_3') \delta_{\sigma_2 \sigma_2'} \delta_{\sigma_3 \sigma_3'} \right).$$

(1)

Here, the matrix element of the current operator between one-body states of the constituent quark coupling to the intermediate boson is taken of the form

$$\langle p_1', \sigma_1' | \hat{J}^{\mu}_{\text{spec}} | p_1, \sigma_1 \rangle = e_1 \bar{u}(p_1', \sigma_1') \gamma^\mu u(p_1, \sigma_1),$$

(2)

for the electromagnetic case, and

$$\langle p_1', \sigma_1' | \hat{J}^{\mu}_{\text{a, spec}} | p_1, \sigma_1 \rangle = \bar{u}(p_1', \sigma_1') \left[g_A \gamma^\mu + \frac{2f_{\rho}}{Q^2 + m_\rho^2} g_{\rho q q} \bar{q}q^\mu \right] \gamma_5 \frac{1}{2} \tau_a u(p_1, \sigma_1),$$

(3)

for the axial case, where $u(p_1, \sigma_1)$ represents the four-component Dirac spinor of quark 1; for more details of the formalism see ref. [11]. Obviously, eq. (2) represents...
the usual relativistic electromagnetic current for a point-like particle, and eq. (3) is the conventional axial current with the pion-pole term included. Now, there are several important features to be noted about the PFSM currents. First, all four components of the momentum transfer $\hat{q}^\mu = p_1^\mu - p_2^\mu$ to the struck quark are in general different from the one of the momentum transfer $q^\mu$ to the baryon as a whole; $q^\mu$ results uniquely from the spectator conditions in eq. (1) and from the conservation of the overall momentum $\mu^\nu = p_1^\nu - p_2^\nu$ (translational invariance). Second, in the PFSM construction of the current in eq. (1), a normalization factor $N$ has to be introduced in order to guarantee for the proper charge normalization (of the proton) [10]. In principle, a Lorentz-invariant form of the normalization factor $N$ involves the interacting masses $M$ and $M'$ of the incoming and outgoing baryon states, respectively, and can be chosen in many different ways [11]

$$N(x, y) = \left( \frac{M}{\sum_i \omega_i} \right)^x \left( \frac{M'}{\sum_i \omega_i} \right)^{x(1-y)},$$

with $0 \leq x$ and $0 \leq y \leq 1$. From the electromagnetic case, the exponent $x$ is fixed to 3 so that the electric form factor $G_E^p$ yields the proper charge of the proton, see the left panel of Figure 4. The exponent $y$ can be constrained specifically by exploiting time-reversal invariance. The latter implies that in the Breit frame the expectation value of the third component $J^{\mu=3}$ of the current operator has to vanish [12]. From the behaviour of $N$ as a function of $y$ in the right panel of Figure 4 one finds that $y = \frac{1}{2}$ meets this constraint (for all values of the momentum transfer $Q^2$). It should be noted that the normalization factor $N$ entering into eq. (1) introduces contributions from the interacting three-quark systems in a non-separable manner and thus makes the PFSM currents effective many-body ones.

In view of the constraints implemented in its construction, the PFSM is found to be specific, since it does not only guarantee for the invariance of the transition amplitudes under the transformations of the whole Poincaré group (including space and time reflections as well as space-time translations) but also allows to fulfill supplementary requirements such as charge normalization. All of these constraints are maintained in any reference frame, because the point-form calculations are performed in a manifestly covariant manner. One may suspect that the relatively good performance of the PFSM approach is due to the fulfilling of these additional conditions beyond Poincaré invariance.

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[1] P. A. M. Dirac, Rev. Mod. Phys. 21, 392 (1949)
[2] W. H. Klink, Phys. Rev. C 58, 3587 (1998)
[3] L. Y. Glozman, W. Plessas, K. Varga and R. F. Wagenbrunn, Phys. Rev. D 58, 094030 (1998)
[4] R. F. Wagenbrunn, S. Boffi, W. Klink, W. Plessas and M. Radici, Phys. Lett. B511, 53 (2001)
[5] L. Y. Glozman, M. Radici, R. F. Wagenbrunn, S. Boffi, W. Klink and W. Plessas, Phys. Lett. B516, 183 (2001)
[6] S. Boffi, L. Y. Glozman, W. Klink, W. Plessas, M. Radici and R. F. Wagenbrunn, Eur. Phys. J. A 14, 17 (2002)
[7] K. Berger, R. F. Wagenbrunn and W. Plessas, Phys. Rev. D 70, 094027 (2004)
[8] R. K. Bhaduri, L. E. Cohler and Y. Yogami, Nuovo Cim. A 65, 376 (1981)
[9] L. Theussl, R. F. Wagenbrunn, B. Desplanques and W. Plessas, Eur. Phys. J. A 12, 91 (2001)
[10] L. Y. Glozman, Z. Papp, W. Plessas, K. Varga and R. F. Wagenbrunn, Phys. Rev. C 57, 3406 (1998)
[11] T. Melde, L. Canton, W. Plessas and R. F. Wagenbrunn, Eur. Phys. J. A 25, 97 (2005)
[12] L. Durand III, P. C. DeCelles and R. B. Marr, Phys. Rev. 126, 1882 (1962)
[13] L. Andivahis et al., Phys. Rev. D 50, 5491 (1994)
[14] R. C. Walker et al., Phys. Lett. B224, 353 (1989)
[15] A. F. Sill et al., Phys. Rev. D 48, 29 (1993)
[16] G. Höhler et al., Nucl. Phys. B114, 505 (1976)
[17] W. Bartel et al., Nucl. Phys. B58, 429 (1973)
[18] T. Eden et al., Phys. Rev. C 50, R1749 (1994)
[19] M. Meyerhoff et al., Phys. Lett. B327, 201 (1994)
[20] A. Lung et al., Phys. Rev. Lett. 70, 718 (1993)
[21] C. Herberg et al., Eur. Phys. J. A 5, 131 (1999)
[22] D. Rohe et al., Phys. Rev. Lett. 83, 4257 (1999)
[23] M. Ostrick et al., Phys. Rev. Lett. 83, 276 (1999)
[24] J. Becker et al., Eur. Phys. J. A 6, 329 (1999); with FSI corrections calculated by J. Golak et al., Phys. Rev. C 63, 034006 (2001)
[25] I. Passchier et al., Phys. Rev. Lett. 82, 4988 (1999)
[26] H. Zhu et al., Phys. Rev. Lett. 87, 081801 (2001)
[27] R. Schiavilla and I. Sick, Phys. Rev. C 64, 041002 (2001)
[28] J. Bermuth et al., Phys. Lett. B564, 199 (2003)
[29] R. Madey et al., Phys. Rev. Lett. 91, 122002 (2003)
FIG. 4: Left: Proton charge as a function of the exponent $x$ in the normalization factor $N$ of eq. (4). Right: Expectation value of the electromagnetic current component $\hat{J}_{\mu=3}$ in the Breit frame as a function of the exponent $y$ in the normalization factor $N$ of eq. (4) for three different values of the momentum transfer $Q^2$.

[30] D. I. Glazier et al., Eur. Phys. J. A 24, 101 (2005)
[31] P. Markowitz et al., Phys. Rev. C 48, R5 (1993)
[32] S. Rock et al., Phys. Rev. Lett. 49, 1139 (1982)
[33] E.E.W. Bruins et al., Phys. Rev. Lett. 75, 21 (1995)
[34] H. Gao et al., Phys. Rev. C 50, R546 (1994)
[35] H. Anklin et al., Phys. Lett. B428, 248 (1998)
[36] H. Anklin et al., Phys. Lett. B336, 313 (1994)
[37] W. Xu et al., Phys. Rev. Lett. 85, 2900 (2000)
[38] G. Kubon et al., Phys. Lett. B524, 26 (2002)
[39] W. Xu et al., Phys. Rev. C 67, 012201 (2003)
[40] It should be emphasized that this normalization factor in the current operator has nothing to do with the normalization of the nucleon states. The necessity of employing $N$ in the PFSM current operator could easily be overlooked. It becomes especially evident, however, when correctly normalized states are used in the evaluation of the matrix element $\hat{m}$, see the details given in ref. [11].