Determination of nucleon form factors from baryonic B decays

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It is the first time that the study of three-body baryonic B decays offers an independent determination of the nucleon form factors for the timelike four momentum transfer ($t > 0$) region, such as $G_M^{p,n}(t)$ of vector currents and those of axial ones. Explicitly, from the data of $\bar{B}^0 \to n\bar{p}D^{*+}$ and $\Lambda\bar{p}\pi^+$ we find a constant ratio of $G_M^n(t)/G_M^p(t) = -1.3 \pm 0.4$, which supports the FENICE experimental result. The vector and axial-vector form factors of $p\bar{n}$, $p\bar{p}$ and $n\bar{n}$ pairs due to weak currents are also presented, which can be tested in future experiments.

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The nucleon form factors still attract attentions as they reveal the hadron structures and play important roles in any scattering or decaying processes involving baryons [1]. These form factors depend on the four momentum transfer (t), which is either spacelike ($t < 0$) or timelike ($t > 0$). The behaviors of the form factors versus t have been extensively studied in various QCD models [2–13], such as perturbative QCD (PQCD), chiral perturbation theory (ChPT), vector meson dominance (VMD) approach, and dispersion relation (DR) method. Experimentally, accurate electromagnetic data for vector form factors with the spacelike momentum transfer have become abundant [14], whereas the data on axial-vector form factors are available only for the spacelike region with $|t| < 1$ GeV^2 in the neutrino-nucleon scattering [1]. In particular, the timelike electromagnetic form factors of the proton have been extracted from $e^+e^- \to p\bar{p}$ ($p\bar{p} \to e^+e^-$) [15], but only a few data points have been collected for those of the neutron by the FENICE Collaboration [16]. Currently, due to experimental difficulties, there are no data on the timelike axial structures, induced from the weak currents due to W and Z bosons. Moreover, there exists some inconsistency between the measurement and theory [2–13] for the $n\bar{n}$ data, unlike the $p\bar{p}$ case, which seems to be well understood by the theoretical calculations. Clearly, more theoretical studies as well as precise experimental measurements on the nucleon form factors are needed to improve our understanding of strong interactions.

In this paper, we shall show that the three-body baryonic B decays of $B \to \mathbf{B}\bar{\mathbf{B}}'M$, such as $\bar{B}^0 \to n\bar{p}D^{*+}$ and $\bar{B}^0 \to \Lambda\bar{p}\pi^+$, can provide valuable information on the nucleon form factors. In general, the three-body baryonic B decays involve timelike form factors from vector, axial-vector, scalar and pseudoscalar currents, respectively. In the scale of $m_b \sim 4$ GeV, the PQCD is suitable for us to systematically examine not only the form factors of vector currents but also those of axial-vector ones. We note that the PQCD approach in a series of works in Refs. [17–23] has been developed as a reliable tool to explain the experimental data on the baryonic B decays.

In the widely used factorization method [24,25], which splits the four quark operators into two currents by the vacuum insertion, there are three types of three-body baryonic B decays: Type I is for the decay in which a meson is transformed from B together with an emitted baryon pair; Type II is for the mode in which a baryon pair is transited from B together with an ejected meson; and Type III is the mixture of Types I and II. With the factorization method, the decay amplitudes for Types I and II are proportional to $\langle \mathbf{B}\bar{\mathbf{B}}'|J_\mu^1|0\rangle\langle M|J_2^\mu|B\rangle$ and $\langle M|J_\mu^1|0\rangle\langle \mathbf{B}\bar{\mathbf{B}}'|J_2^\mu|B\rangle$, respectively. For the present measured modes, for instance, $\bar{B}^0 \to n\bar{p}D^{*+}$ [26] and $\bar{B}^0 \to \Lambda\bar{p}\pi^+$ [27] belong to Type I, while $\bar{B}^0 \to p\bar{p}D^{(*)0}$ [28] and $B^- \to \Lambda\bar{p}J/\Psi$ [29] are classified as Type II, whereas $B^- \to p\bar{p}K^{(*)-}$, $\bar{B}^0 \to p\bar{p}K_S$, $B^- \to p\bar{p}\pi^-$ [30] and $B^- \to \Lambda\bar{\Lambda}K^-$ [31] are of Type III. Although the decay modes of Types I and III are of our current interest, those of Type III are inevitably affected by the uncertainties of the $B \to \mathbf{B}\bar{\mathbf{B}}'$ transition form factors. In this study, we shall concentrate on the Type I modes of $\bar{B}^0 \to n\bar{p}D^{*+}$ and $\bar{B}^0 \to \Lambda\bar{p}\pi^+$.

With the effective Hamiltonians [32] at the quark level, the decay amplitudes are given by [18–22]

$$\mathcal{A}(\bar{B}^0 \to n\bar{p}D^{*+}) = \frac{G_F}{\sqrt{2}} V_{cb}V_{ud}^* a_1 \langle n\bar{p}|(\bar{d}u)_{V-A}|0\rangle \times \langle D^{*+}|(\bar{c}b)_{V-A}|\bar{B}^0\rangle, \tag{1}$$

$$\begin{aligned}\mathcal{A}(\bar{B}^0 \to \Lambda\bar{p}\pi^+) = \frac{G_F}{\sqrt{2}}\{&(V_{ub}V_{us}^* a_1 - V_{tb}V_{ts}^* a_4) \\ &\times \langle \Lambda\bar{p}|(\bar{s}u)_{V-A}|0\rangle\langle \pi^+|(\bar{u}b)_{V-A}|\bar{B}^0\rangle \\ &+ V_{tb}V_{ts}^* 2a_6 \langle \Lambda\bar{p}|(\bar{s}u)_{S+P}|0\rangle \\ &\times \langle \pi^+|(\bar{u}b)_{S-P}|\bar{B}^0\rangle\},\end{aligned} \tag{2}$$

where G_F is the Fermi constant, $V_{q_iq_j}$ are the CKM matrix elements, $(q_iq_j)_{V-A} = q_i\gamma_\mu(1-\gamma_5)q_j$, $(q_iq_j)_{S\pm P} = q_i(1\pm\gamma_5)q_j$ and a_i ($i = 1, 4, 6$) are given by

$$a_1 = c_1^{\rm eff} + \frac{1}{N_c^{\rm eff}} c_2^{\rm eff}, \qquad a_4 = c_4^{\rm eff} + \frac{1}{N_c^{\rm eff}} c_3^{\rm eff}, \qquad a_6 = c_6^{\rm eff} + \frac{1}{N_c^{\rm eff}} c_5^{\rm eff}, \tag{3}$$

with $c_i^{\rm eff}$ $(i = 1, 2, \cdots, 6)$ being the effective Wilson coefficients (WCs) shown in Refs. [32] and $N_c^{\rm eff}$ the effective color number. Here, we have used the generalized factorization method with the nonfactorizable effect absorbed in $N_c^{\rm eff}$. For the matrix elements of $\langle \pi^+ | (\bar{u}b)_{V-A} | \bar{B}^0 \rangle$ and $\langle D^{*+} | (\bar{c}b)_{V-A} | \bar{B}^0 \rangle$, we use the results in Refs. [25,33]. For the timelike baryonic form factors, we have

$$\begin{aligned}
\langle \mathbf{B}\bar{\mathbf{B}}' | \bar{q}_i \gamma_\mu q_j | 0 \rangle &= \bar{u}(p_{\mathbf{B}}) \left\{ F_1(t)\gamma_\mu + \frac{F_2(t)}{m_{\mathbf{B}} + m_{\bar{\mathbf{B}}'}} i\sigma_{\mu\nu}(p_{\bar{\mathbf{B}}'} + p_{\mathbf{B}})_\mu \right\} v(p_{\bar{\mathbf{B}}'}) \\
&= \bar{u}(p_{\mathbf{B}}) \left\{ [F_1(t) + F_2(t)]\gamma_\mu + \frac{F_2(t)}{m_{\mathbf{B}} + m_{\bar{\mathbf{B}}}} (p_{\bar{\mathbf{B}}} - p_{\mathbf{B}})_\mu \right\} v(p_{\bar{\mathbf{B}}}), \\
\langle \mathbf{B}\bar{\mathbf{B}}' | \bar{q}_i \gamma_\mu \gamma_5 q_j | 0 \rangle &= \bar{u}(p_{\mathbf{B}}) \left\{ g_A(t)\gamma_\mu + \frac{h_A(t)}{m_{\mathbf{B}} + m_{\bar{\mathbf{B}}'}} (p_{\bar{\mathbf{B}}'} + p_{\mathbf{B}})_\mu \right\} \gamma_5 v(p_{\bar{\mathbf{B}}}), \\
\langle \mathbf{B}\bar{\mathbf{B}}' | \bar{q}_i q_j | 0 \rangle &= f_S(t) \bar{u}(p_{\mathbf{B}}) v(p_{\bar{\mathbf{B}}'}), \\
\langle \mathbf{B}\bar{\mathbf{B}}' | \bar{q}_i \gamma_5 q_j | 0 \rangle &= g_P(t) \bar{u}(p_{\mathbf{B}}) \gamma_5 v(p_{\bar{\mathbf{B}}'}),
\end{aligned} \tag{4}$$

where the four momentum transfer in the timelike region is $t = (p_{\mathbf{B}} + p_{\bar{\mathbf{B}}'})^2$, $q_i = u$, d, and s, and F_1, F_2, g_A, h_A, f_S, and g_P are the form factors.

In this paper, we will study the form factors in Eq. (4) based on the experimental data in the baryonic B decays. We begin by defining the baryonic form factors of F_1 and g_A by

$$\langle \mathbf{B} | J_\mu^{\rm em} | \mathbf{B}' \rangle = \bar{u}(p_{\mathbf{B}})[F_1(t)\gamma_\mu + g_A(t)\gamma_\mu\gamma_5] v(p_{\mathbf{B}'}), \tag{5}$$

where $J_\mu^{\rm em} = Q_q \bar{q}\gamma_\mu q$ and $t = (p_{\mathbf{B}} - p_{\mathbf{B}'})^2$. Note that F_2 and h_A are not included in Eq. (5) due to the helicity conservation. To exhibit the chirality or helicity, we rewrite Eq. (5) as

$$\langle \mathbf{B}_{\uparrow+\downarrow} | J_\mu^{\rm em} | \mathbf{B}'_{\uparrow+\downarrow} \rangle = \bar{u}(p_{\mathbf{B}}) \left[\gamma_\mu \frac{1+\gamma_5}{2} G^\uparrow(t) + \gamma_\mu \frac{1-\gamma_5}{2} G^\downarrow(t) \right] u(p_{\mathbf{B}'}), \tag{6}$$

where $|\mathbf{B}_{\uparrow+\downarrow}\rangle \equiv |\mathbf{B}_\uparrow\rangle + |\mathbf{B}_\downarrow\rangle$ respects both flavor $SU(3)$ and spin $SU(2)$ symmetries, e.g., $|p_\uparrow\rangle = \sqrt{1/18}[u_\uparrow u_\downarrow d_\uparrow + u_\downarrow u_\uparrow d_\uparrow - 2u_\uparrow u_\uparrow d_\downarrow + \text{permutations}]$. In Eq. (6), $G^\uparrow(t)$ and $G^\downarrow(t)$ represent the right-handed and left-handed form factors, which can be further decomposed as

$$\begin{aligned}
G^\uparrow(t) &= e_\parallel^\uparrow G_\parallel(t) + e_{\bar\parallel}^\uparrow G_{\bar\parallel}(t), \\
G^\downarrow(t) &= e_\parallel^\downarrow G_\parallel(t) + e_{\bar\parallel}^\downarrow G_{\bar\parallel}(t),
\end{aligned} \tag{7}$$

where the constants $e_\parallel^{\uparrow(\downarrow)}$ and $e_{\bar\parallel}^{\uparrow(\downarrow)}$ are defined by

$$e_\parallel^{\uparrow(\downarrow)} = \langle \mathbf{B}_{\uparrow(\downarrow)} | \mathbf{Q}_\parallel | \mathbf{B}'_{\uparrow(\downarrow)} \rangle, \qquad e_{\bar\parallel}^{\uparrow(\downarrow)} = \langle \mathbf{B}_{\uparrow(\downarrow)} | \mathbf{Q}_{\bar\parallel} | \mathbf{B}'_{\uparrow(\downarrow)} \rangle, \tag{8}$$

respectively, with $\mathbf{Q}_{\parallel(\bar\parallel)} = \sum_i Q_{\parallel(\bar\parallel)}(i)$. In Eq. (8), the summation is over the charges carried by the valence quarks $(i = 1, 2, 3)$ in the baryon with helicities parallel ($\parallel$) and antiparallel ($\bar\parallel$) to the baryon spin directions of $(\uparrow, \downarrow)$. Since $G_{\parallel(\bar\parallel)}(t)$ are the form factors accompanied by the (anti-)parallel hard-scattering amplitudes with baryon wave functions, based on the QCD counting rules in the PQCD [2,34,35], they can be expressed by

$$G_\parallel(t) = \frac{C_\parallel}{t^2} \left[\ln\left(\frac{t}{\Lambda_0^2}\right) \right]^{-\gamma}, \qquad G_{\bar\parallel}(t) = \frac{C_{\bar\parallel}}{t^2} \left[\ln\left(\frac{t}{\Lambda_0^2}\right) \right]^{-\gamma}, \tag{9}$$

where $\gamma = 2.148$, $\Lambda_0 = 300$ MeV [6] and $C_{\parallel,\bar\parallel}$ are parameters to be determined. We remark that the asymptotic formulas in Eq. (9) are exact only when t is large. For smaller t, such as when t being close to the two-nucleon threshold, higher power corrections are expected [36]. The corrections will be averaged into the errors of $C_{\parallel,\bar\parallel}$ in our data fitting. From Eqs. (5)–(8), we have

$$\begin{aligned}
F_1(t) &= (e_\parallel^\uparrow + e_\parallel^\downarrow) G_\parallel(t) + (e_{\bar\parallel}^\uparrow + e_{\bar\parallel}^\downarrow) G_{\bar\parallel}(t), \\
g_A(t) &= (e_\parallel^\uparrow - e_\parallel^\downarrow) G_\parallel(t) + (e_{\bar\parallel}^\uparrow - e_{\bar\parallel}^\downarrow) G_{\bar\parallel}(t).
\end{aligned} \tag{10}$$

Although the above equations are derived in the spacelike region, the timelike form factors can be easily written via the crossing symmetry [6,37], which transforms the particle in the initial state to its antiparticle in the final state and reverses its helicity. However, in general, the values of the timelike $G_\parallel(t)$ and $G_{\bar\parallel}(t)$ form factors are complex numbers unlike the spacelike ones for which there is a time reversal symmetry between initial and final states.

In Eq. (4), F_2 is suppressed by $1/(t\ln[t/\Lambda_0^2])$ in comparison with F_1 [38,39] and therefore can be safely ignored, while g_P is found to be related to f_S as [20]

$$g_P = f_S, \tag{11}$$

and f_S and h_A are deduced from equation of motion, given by

$$f_S(t) = \frac{m_{\mathbf{B}} - m_{\mathbf{B}'}}{m_{q_i} - m_{q_j}} F_1(t), \qquad h_A(t) = -\frac{(m_{\mathbf{B}} + m_{\mathbf{B}'})^2}{t} g_A(t). \tag{12}$$

Thus, once we figure out F_1 and g_A, all other form factors f_S, h_A and g_P will be determined in terms of Eqs. (11) and (12). For the electromagnetic current $J_\mu^{\rm em} = \frac{2}{3}\bar{u}\gamma_\mu u - \frac{1}{3}\bar{d}\gamma_\mu d$, from Eq. (10) it is clear that $g_A = 0$ since $e^{\uparrow}_{\parallel(\bar{\parallel})} = e^{\downarrow}_{\parallel(\bar{\parallel})}$. Furthermore, from Eqs. (4)–(8) we have

$$G_M^p(t) \equiv F_1^{p\bar{p}}(t) + F_2^{p\bar{p}}(t) \simeq F_{1({\rm em})}^{p\bar{p}}(t) = G_\parallel(t),$$
$$G_M^n(t) \equiv F_1^{n\bar{n}}(t) + F_2^{n\bar{n}}(t) \simeq F_{1({\rm em})}^{n\bar{n}}(t) = -\frac{G_\parallel(t)}{3} + \frac{G_{\bar{\parallel}}(t)}{3}, \tag{13}$$

where we have neglected the small $F_2^{N\bar{N}}$ terms comparing with those of $F_1^{N\bar{N}}$ [38,39]. Similarly, for the Z coupled current $J_\mu^Z = \frac{1}{2}\bar{u}\gamma_\mu(\frac{1-\gamma_5}{2})u - \frac{1}{2}\bar{d}\gamma_\mu(\frac{1-\gamma_5}{2})d - \sin^2\theta_W J_\mu^{\rm em}$, we use the same form factors defined in Eqs. (5) and (6) by replacing $J_\mu^{\rm em}$ with J_μ^Z and we get

$$F_{1(Z)}^{p\bar{p}}(t) = \frac{2 - 3\sin^2\theta_W}{3} G_\parallel(t) - \frac{1}{6} G_{\bar{\parallel}}(t),$$
$$g_{A(Z)}^{p\bar{p}}(t) = \frac{2}{3} G_\parallel(t) + \frac{1}{6} G_{\bar{\parallel}}(t),$$
$$F_{1(Z)}^{n\bar{n}}(t) = \frac{-2 - \sin^2\theta_W}{3} G_\parallel(t) + \frac{1 + 2\sin^2\theta_W}{6} G_{\bar{\parallel}}(t),$$
$$g_{A(Z)}^{n\bar{n}}(t) = -\frac{2}{3} G_\parallel(t) - \frac{1}{6} G_{\bar{\parallel}}(t). \tag{14}$$

Since the behaviors of $G_\parallel(t)$ and $G_{\bar{\parallel}}(t)$ have been given in Eq. (9), what we shall do next is to fix the parameters $C_\parallel$ and $C_{\bar{\parallel}}$ in terms of

$$F_1^{n\bar{p}}(t) = \tfrac{4}{3}G_\parallel(t) - \tfrac{1}{3}G_{\bar{\parallel}}(t), \qquad g_A^{n\bar{p}}(t) = \tfrac{4}{3}G_\parallel(t) + \tfrac{1}{3}G_{\bar{\parallel}}(t),$$
$$F_1^{\Lambda\bar{p}}(t) = \sqrt{\tfrac{3}{2}}G_\parallel(t), \qquad g_A^{\Lambda\bar{p}}(t) = \sqrt{\tfrac{3}{2}}G_\parallel(t), \tag{15}$$

derived from Eqs. (1)–(8) with the data in $\bar{B}^0 \to n\bar{p}D^{*+}$ and $\bar{B}^0 \to \Lambda\bar{p}\pi^+$.

Before performing the numerical analysis, we would like to briefly discuss the generalized factorization method. It is known that the factorization method [24,25] suffers from several possible hadron uncertainties, such as those from the nonfactorizable effect, annihilation contribution, and final state interaction. To describe these uncertainties, we take the decay of $\bar{B}^0 \to n\bar{p}D^{*+}$ as an example, while those for $\bar{B}^0 \to \Lambda\bar{p}\pi^+$ can be treated in a similar manner. The amplitude of $\bar{B}^0 \to n\bar{p}D^{*+}$ from the color suppressed operator is given by

$$c_2^{\rm eff}\langle n\bar{p}D^{*+}|(\bar{d}_\alpha u_\beta)_{V-A}(\bar{c}_\beta b_\alpha)_{V-A}|\bar{B}^0\rangle = \frac{c_2^{\rm eff}}{N_c}\langle n\bar{p}|(\bar{d}u)_{V-A}|0\rangle\langle D^{*+}|(\bar{c}b)_{V-A}|\bar{B}^0\rangle + \frac{c_2^{\rm eff}}{2} \times \langle n\bar{p}D^{*+}|(\bar{d}\lambda^a u)_{V-A}(\bar{c}\lambda^a b)_{V-A}|\bar{B}^0\rangle, \tag{16}$$

where $\delta_{\beta\beta'}\delta_{\alpha\alpha'} = \delta_{\beta\alpha}\delta_{\alpha'\beta'}/N_c + \lambda^a_{\beta\alpha}\lambda^a_{\alpha'\beta'}/2$ has been used to deal with color index α (β) and the second term on the right-hand side is the so-called nonfactorizable effect. Although the nonfactorizable effect cannot be directly and unambiguously determined by theoretical calculations, in the generalized factorization method [32], this contribution can be absorbed in the effective color number $N_c^{\rm eff}$ running from 2 to ∞ in Eq. (3). The amplitude of the annihilation contribution is given by

$$\mathcal{A}_{\rm an}(\bar{B}^0 \to n\bar{p}D^{*+}) = \frac{G_F}{\sqrt{2}} V_{cb}V_{ud}^* a_2 \langle n\bar{p}D^{*+}|(\bar{c}u)_{V-A}|0\rangle \times \langle 0|(\bar{d}b)_{V-A}|\bar{B}^0\rangle, \tag{17}$$

where a_2 is around (0.1–0.01)a_1, which is suppressed. In addition, based on the power expansion of $1/t$ in the PQCD approach, $\mathcal{A}_{\rm an}(\bar{B}^0 \to n\bar{p}D^{*+}) \propto 1/t^3$, which is much suppressed than $\mathcal{A}(\bar{B}^0 \to n\bar{p}D^{*+})$ as $t \sim m_b^2$ [40]. For the final state interaction, the most possible source is via the two-particle rescattering to the baryon pair, such as $\bar{B}^0 \to M_1M_2D^{*+} \to n\bar{p}D^{*+}$ with $M_{1,2}$ representing meson states. However, such processes would shape the curve associated with the phase spaces in the decay rate distributions [41–43], which have been excluded in the charmless baryonic B decay experiments.

In our numerical analysis, we take $c_i^{\rm eff}(i = 1, 2, \cdots, 6) \simeq (1.17, -0.37, 0.0246, -0.0523, 0.0154, -0.066)$ [32,44,45], $m_u(m_b) = 3.2$ MeV, $m_s(m_b) = 90$ MeV [32] and $m_b(m_b) = 4.2$ GeV [46], and the weak phase $\gamma = 59.8° \pm 4.9°$ [47]. For the experimental data in the B decays, we use [26,27]

$${\rm Br}(\bar{B}^0 \to n\bar{p}D^{*+}) = (14.5 \pm 4.3) \times 10^{-4},$$
$${\rm Br}(\bar{B}^0 \to \Lambda\bar{p}\pi^+) = (3.29 \pm 0.47) \times 10^{-6}, \tag{18}$$

and the other available data in Ref. [27], such as spectrum vs invariant mass and angular distributions.

Based on the χ^2 fitting, we obtain

$$\chi^2/{\rm dof} = 0.7, \qquad 1/N_c^{\rm eff} = 0.2 \pm 0.2, \tag{19}$$

where dof denotes the degree of freedom. It is clear that $\chi^2/{\rm dof} = 0.7$ presents a reliable fit, while $N_c^{\rm eff} \sim 2.5$–∞ means a limited nonfactorizable effects with the small contributions from the annihilation and final state interaction. We stress that if the hadronic uncertainties were large, $N_c^{\rm eff}$ from 2 to ∞ would not be accounted for in the data. The coefficients of $C_\parallel$ and $C_{\bar{\parallel}}$ in Eq. (9) are found to be

$$C_{\parallel} = 83.7 \pm 5.7\ \mathrm{GeV}^4 \quad \text{and} \quad C_{\bar{\parallel}} = -246.3 \pm 92.1\ \mathrm{GeV}^4 \tag{20}$$

which lead to

$$G_{\bar{\parallel}}(t)/G_{\parallel}(t) = -2.9 \pm 1.1. \tag{21}$$

Here, we have assumed that both $C_{\parallel}$ and $C_{\bar{\parallel}}$ are real since their imaginary parts are expected to be small based on the argument in Refs. [12,37] as well as the result in the DR method [48].

As seen in Fig. 1(a), the fitted values of $G_M^n(t)$ extracted from the B decays are in agreement with the FENICE data [16] by the assumptions of $|G_M^n| = |G_E^n|$ and $|G_E^n| = 0$. We note that $G_E^n(t) \equiv F_1^{n\bar{n}}(t) + \frac{t}{4M_n^2} F_2^{n\bar{n}}$ is the neutron electric form factor. Clearly, in our calculation based on the QCD counting rule, $G_E^n \sim G_M^n$.

From Eqs. (9), (13), and (20), we get

$$G_M^n(t)/G_M^p(t) = -1.3 \pm 0.4, \tag{22}$$

which supports the measurements in Refs. [15,16]. We note that in Eq. (22) the ratio is a constant due to the same power expansions in Eq. (9) and the minus sign is necessary in order to match the measured baryonic B decay branching ratios, which also confirms the theoretical result in Refs. [2,3,34,35]. Moreover, the puzzle of $|G_M^n(t)/G_M^p(t)| \geq 1$ is also solved as indicated in Eq. (22). We note that the ratio $G_M^n(t)/G_M^p(t)$ is predicted to be $-2/3$ and $-1/2$ in the QCD counting [2] and sum rules [3], respectively, while the DR [4–7] and VMD [8,9] methods yield only half of the values indicated by the data points. It is interesting to see that if $G_{\parallel}(t)/G_{\bar{\parallel}}(t) = -1$ instead of the fitted one in Eq. (21), from Eq. (13) $G_P^M(t)/G_M^p(t) = -2/3$ is recovered as in Ref. [2]. Therefore, our result on the ratio in Eq. (21) could help us to improve the PQCD calculations.

Since we have related all the nucleon form factors as in Eqs. (13)–(15), we find

$$\begin{aligned}
F_1^{n\bar{p}}(t) &= (193.7 \pm 31.6) G_p(t),\\
g_A^{n\bar{p}}(t) &= (29.5 \pm 31.6) G_p(t),\\
F_{1(Z)}^{p\bar{p}}(t) &= (78.4 \pm 15.6) G_p(t),\\
g_{A(Z)}^{p\bar{p}}(t) &= (14.8 \pm 15.8) G_p(t),\\
F_{1(Z)}^{n\bar{n}}(t) &= (-121.1 \pm 22.5) G_p(t),\\
g_{A(Z)}^{n\bar{n}}(t) &= (-14.8 \pm 15.8) G_p(t),
\end{aligned} \tag{23}$$

where $G_p(t) \equiv 1/t^2[\ln(t/\Lambda_0^2)]^{-\gamma}$. The central values of $F_1^{N\bar{N}'}$ and $g_A^{N\bar{N}'}$ in Eq. (23) are shown in Figs. 1(b) and 1(c), respectively. The valid ranges for these timelike form factors are the same as those in the three-body baryonic B decays of $B \to \mathbf{B}\bar{\mathbf{B}}'M$, i.e., $(m_{\mathbf{B}} + m_{\mathbf{B}'})^2 \simeq 4\ \mathrm{GeV}^2 \leq t \leq (m_B - m_M)^2 \simeq 16–25\ \mathrm{GeV}^2$.

In sum, we have shown that the study of the measured three-body baryonic B decays of $\bar{B}^0 \to n\bar{p}D^{*+}$ and $\bar{B}^0 \to \Lambda\bar{p}\pi^+$ leads to $G_P^n(t)/G_t^p(t) = -1.3 \pm 0.4$, which supports the FENICE measurement. The minus sign for the ratio $G_P^n(t)/G_t^p(t)$, given by the previous theoretical calculations, has been enforced to fit the B decay data. We have pointed out that our fitted value for the ratio of $G_{\parallel}(t)$ and $G_{\bar{\parallel}}(t)$ may be useful for us to perform various QCD calculations on $e^+e^- \to n\bar{n}$. We have also predicted the timelike vector and axial-vector nucleon form factors induced from the weak currents, such as $F_1^{N\bar{N}'}(t)$ and $g_A^{N\bar{N}'}(t)$ ($N, N' = p$ and n).

Finally, we remark that apart from the use of $\bar{B}^0 \to p\bar{n}D^{*+}$ and $\bar{B}^0 \to \Lambda\bar{p}\pi^+$, there are more decays directly connecting to the timelike form factors, such as $\bar{B}^0 \to p\bar{n}(D^+, \rho^+, \pi^+)$ and $\bar{B}^0 \to \Lambda\bar{p}\rho^+$ as well as the corresponding charged B modes, which are within the accessibility of the current B factories at KEK and SLAC. It is clear that as more and more data becomes available from current and future B factories, the nucleon form factors can be further constrained and determined. Moreover, the new measurements in $e^+e^- \to n\bar{n}$ are progressing in DAΦNE at Frascati [8] and planning in PANDA and PAX at GSI [7]. As for the weak nucleon form factors, since the scattering

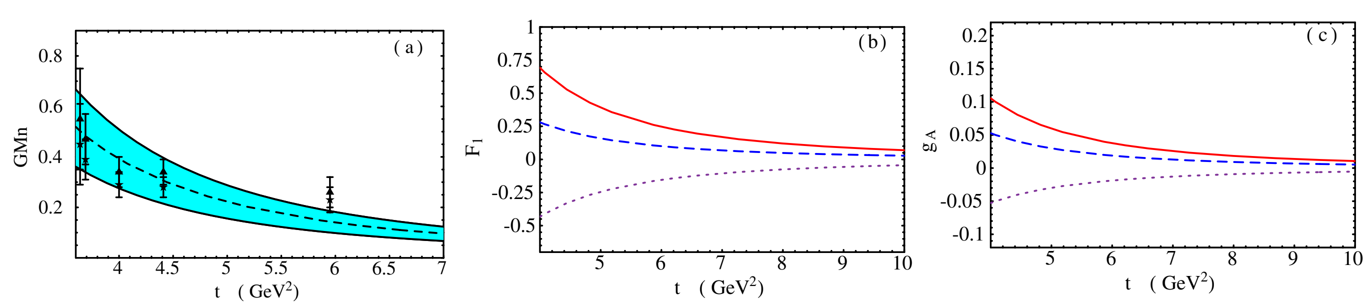

FIG. 1 (color online). Form factors of (a) $G_M^n(t)$, (b) $F_1(t)$, and (c) $g_A(t)$ with the timelike four momentum transfer t, where the star and triangle symbols represent the FENICE data [16] with the assumptions of $|G_M^n| = |G_E^n|$ and $|G_E^n| = 0$, and the solid, dash, and dotted curve stand for (b) $F_1^{p\bar{n}}(t)$, $F_{1(Z)}^{p\bar{p}}(t)$, and $F_{1(Z)}^{n\bar{n}}(t)$ and (c) $g_A^{p\bar{n}}(t)$, $g_{A(Z)}^{p\bar{p}}(t)$, and $g_{A(Z)}^{n\bar{n}}(t)$, respectively.

of e^+e^- at *BABAR* is at the m_B scale, $F_1^{N\bar{N}'}(t)$ and $g_A^{N\bar{N}'}(t)$ ($N, N' = p$ and n) can be studied via the left-right helicity asymmetry [36] of $A_{PV} = (d\sigma_R - d\sigma_L)/(d\sigma_R + d\sigma_L)$ as in the SAMPLE experiment [1].

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