Pramana – J. Phys., Vol. 37, No. 1, July 1991, pp. 47–56. © Printed in India.

Weak electric and magnetic form factors for semileptonic baryon decays in a relativistic quark model

S N JENA and S PANDA
Department of Physics, Berhampur University, Berhampur 760 007, India

MS received 3 December 1990

Abstract. Weak electric and magnetic form factors for semileptonic baryon decays are calculated in a relativistic quark model based on the Dirac equation with the independent-quark confining potential of the form $V_q(r) = \frac{1}{2}(1+\gamma^0)(a^2 r + V_0)$. The values obtained for (g_2/g_1) are not very much different from the nonrelativistic results of Donoghue and Holstein. The values of (g_1/f_1) extracted from our model calculations of (f_2/f_1) in the Cabibbo limit compare well with the experimental values. The values of (f_2/f_1) for various semileptonic transitions are also estimated incorporating phenomenologically the effect of nonzero g_2 in the ratio (g_1/f_1). It is found that the SU(3)-symmetry breaking does not generate significant departures in (f_2/f_1) values from the corresponding Cabibbo predictions.

Keywords. Weak electric and magnetic form factors; SU(3)-symmetry breaking; quark model; semileptonic baryon decays; Cabibbo fits.

PACS Nos 12·40; 13·40

1. Introduction

In our recent studies (Jena and Panda 1989, 1990), we have investigated a relativistic independent-quark model of the form

$$V_q(r) = \tfrac{1}{2}(1+\gamma^0)(a^2 r + V_0), \quad a > 0 \tag{1}$$

for the baryon core, where the quarks inside the core are assumed to move relatively independently in a phenomenologically effective confining potential in the zeroth order. Such a potential model describing the quark dynamics inside the core has been found to be quite suitable in predicting the centre-of-mass motion corrected values for the static properties of S-wave baryons within a reasonable limit of agreement with the experimental data. In the same work we have shown the model estimation of the axial vector coupling constants for weak β-decays of octet baryons. In the present work we would like to extend such investigation further in the context of the Cabibbo analysis (Bohm *et al* 1984; Garcia and Kielanowski 1982 and Donoghue and Holstein 1982) to find a reasonable estimation of the possible SU(3) symmetry breaking effects in semileptonic baryon decays in terms of the weak electric and magnetic form factors.

Semileptonic baryon decays of spin-$\frac{1}{2}$ octet baryons have provided a very successful ground for testing the Cabibbo theory (Cabibbo 1963) of weak interaction. The interaction Hamiltonian for the semileptonic decay process is given by

$$H_I = \frac{G}{\sqrt{2}} j_\mu l^\mu + h{\cdot}C \tag{2}$$

where $G \simeq 1{\cdot}7 \times 10^{-5}\,\mathrm{GeV}^{-2}$ is the Fermi coupling constant.

$$l^\mu = \bar{\psi}_e \gamma^\mu (1 - \gamma_5) \psi_{\nu_e} \tag{3}$$

is the lepton current and

$$j_\mu = \cos\theta_c \bar{\psi}_\mu \gamma_\mu (1 - \gamma_5) \psi_d + \sin\theta_c \bar{\psi}_\mu \gamma_\mu (1 - \gamma_5) \psi_d \tag{4}$$

is the baryon current with θ_c as the Cabibbo angle. The transition matrix element for the decay of an initial baryon B_1 to final baryon B_2 and a lepton l with its antineutrino (i.e. $B_1 \to B_2 l \bar{\nu}_1$) is proportional to the matrix element of the baryonic weak current which is conventionally expressed in terms of weak form factors $f_i(q^2)$ and $g_i(q^2)$ with $i = 1, 2, 3$ as follows:

$$\langle B_2 \uparrow | j^\mu(x) | B_1 \uparrow \rangle = \langle B_2 \uparrow | V^\mu(x) | B_1 \uparrow \rangle + \langle B_2 \uparrow | A^\mu(x) | B_1 \uparrow \rangle \tag{5}$$

where

$$\langle B_2 \uparrow | V^\mu(x) | B_1 \uparrow \rangle = C \exp(iqx) \bar{u}_2(p_2) \left[\gamma^\mu f_1(q^2) + i\sigma^{\mu\nu} q_\nu \frac{f_2(q^2)}{M_1 + M_2} + \frac{q^\mu f_3(q^2)}{M_1 + M_2} \right] u_1(p_1) \tag{6}$$

and

$$\langle B_2 \uparrow | A^\mu(x) | B_1 \uparrow \rangle = C \exp(iqx) \bar{u}_2(p_2) \left[\gamma^\mu g_1(q^2) + i\sigma^{\mu\nu} q_\nu \frac{g_2(q^2)}{M_1 + M_2} + \frac{q^\mu g_3(q^2)}{M_1 + M_2} \right] \gamma^5 u_1(p_1). \tag{7}$$

Here $C = \cos\theta_c (C = \sin\theta_c)$ for $\Delta S = 0 (\Delta S = 1)$ transitions, p_1 and p_2 are the 4-momenta and M_1, M_2 are the masses of initial and final baryons respectively. f_1 and g_1 respectively are the vector and axial vector form factors, f_2 and g_2 respectively are the weak magnetic and weak electric form factors and f_3 and g_3 respectively represent the induced scalar and pseudoscalar form factors. Considering the electronic decay modes in particular the momentum transfer $q^2 = (p_1 - p_2)^2$ is never larger, for which the q^2-dependence of the weak form factors $f_i(q^2)$ and $g_i(q^2)$ is usually neglected for such decays.

The values of f_i and g_i are subject to restrictions imposed by (i) time reversal invariance (ii) conserved vector current hypothesis (CVC) (iii) Goldberger Treimann (GT) relation and (iv) the G-parity invariance. The time reversal invariance implies that the f_i and g_i are real. According to CVC hypothesis the weak vector and electromagnetic current are members of the same SU(3) octet. Thus, f_2 is determined by the anomalous magnetic moments of the proton and the neutron. In the limit of the exact SU(3) symmetry, the absence of second class currents leads to $f_3 = 0 = g_2$. The vanishing of f_3 and g_2 is also a consequence of G-parity invariance. Again the contribution of f_3 and g_3 terms are proportional to m_l^2 as a result of which they are usually assumed to be negligible for electronic decay modes.

It is to be noted that the value of g_2 is zero if SU(3) symmetry is exact. But SU(3) is a badly broken symmetry in the real world. One, therefore, cannot expect strict zero values for the weak electric form factor g_2. While SU(3) breaking effects are expected to contribute to second order in relative mass differences for the vector parts of the hadronic weak current, such contributions are in first order for the axial vector part. Therefore, one anticipates a g_2 term such that $(g_2/g_1) \sim (\Delta M_B/M_B)$. A crude estimate (Donoghue and Holstein 1982) in this manner provides

$$|g_2|_{n \to pe^- \bar{\nu}_e} \lesssim 0{\cdot}03$$

and

$$|g_2|_{\Lambda \to pe^- \bar{\nu}_e} \lesssim 0{\cdot}3$$

which suggests that although g_2 can be neglected in the analysis of neutron beta-decay, it may be relevant in $\Delta S = 1$ semileptonic hyperon decay. There can be a few consequences which the nonzero g_2 term would have in a Cabibbo fit. First of all, the calculated decay rates are expected to be sensitive to nonvanishing g_2 mainly through the interference term $\mathrm{Re}(g_1 g_2^*)$. Secondly since the experimental results obtained for g_1/f_1 in a Cabibbo fit depend on the assumed values of g_2, they may be modified. In view of this it is necessary to know exactly how large is g_2 from various model calculations.

There have been some earlier investigations on this question by several authors (Garcia 1971; Calaprice *et al* 1975; Sugimoto *et al* 1975). On the basis of the existing data available then for $\Lambda \to pe\bar{\nu}_e$, a best fit was obtained by Garcia (1971) with $(g_2/f_1) = -5$. Measurement of correlation between nuclear spin and electron momentum similarly determined a large value of $(g_2/f_1) = -8$ in nuclear beta decay (Calaprice *et al* 1975). However, later studies (Tribble and May 1978; Keppinger *et al* 1978; Lebrun *et al* 1978; Brandle *et al* 1978) show considerable doubts on these results. More recent investigations in some quark model approaches (Donoghus and Holstein 1982; Kobudera *et al* 1985; Eich *et al* 1985; Barik *et al* 1985; Rath and Jena 1989) provide the values of the weak form factors for various semileptonic transitions. The results obtained in all these approaches are more or less identical. Although a slight improvement in the Cabibbo fit (Bourguin *et al* 1982, 1983) has been noticed by including g_2-values of Donoghue and Holstein (1982), it is generally observed that the experimental data are not sufficiently precise to establish the presence of the estimated non-zero g_2. Therefore, for the sake of providing bench marks for future experimental comparisons, it may be worthwhile to calculate the weak form factors for various semileptonic decays from different quark models.

The present work is an attempt in this direction in the same light as that of the earlier ones (Donoghue and Holstein 1982; Barik *et al* 1985 and Rath and Jena 1989) but with a different form of the potential as given in (1). The present potential form has been successfully used by us not only for calculating the magnetic moments of light, charmed and b-flavoured baryons (Jena and Panda 1990a) but also for studying the electromagnetic form factors of the nucleons in the low-q^2 region (Jena and Panda 1990b). In view of this success, it is worthwhile to extend the applicability of the model in the present work to the computation of the weak electric and magnetic form factors in semileptonic baryon transitions.

The present work is organized as follows. In § 2, we briefly show how the weak electric and magnetic form factors are related to the quark wave-function integrals and the physical baryon masses. The quark wave-function integrals appearing in the expressions for weak form factors are evaluated in the present potential model in § 3. Using the model parameters from our earlier work (Jena and Panda 1990b) we present in § 4 the calculated results for weak electric and magnetic form factors corresponding to various semileptonic baryon decays.

2. Weak form factors

A reasonable and reliable estimate of the SU(3) breaking effect in generating a nonzero weak electric form factor g_2 in semileptonic baryon decays can be provided by a relativistic independent quark model describing the baryons. In such models one assumes that the quarks in a baryon core move independently in an average flavour-independent confining potential $V_q(r)$ obeying the Dirac equation, so that a solution to the normalized independent quark wave function $\psi_q(\mathbf{r})$ corresponding to the ground state baryons in the nucleon octet can be written in the form

$$\psi_q = \frac{N_q}{\sqrt{4\pi}} \begin{bmatrix} ig_q(r)/r \\ \boldsymbol{\sigma}\cdot\hat{r} f_q(r)/r \end{bmatrix} \chi\uparrow. \tag{8}$$

Since the independent motion of quarks inside the baryon core does not lead to a state of definite total momentum, as it should to represent the physical state of a baryon, one can use the wave packet formalism of Donoghue and Johnson (1980) to decompose the static three quark baryon core state into components of plane-wave momentum eigenstates.

Interpreting the weak beta decays $B_1 \to B_2 e^- \bar{\nu}_e$ as the quark beta decays $q_i \to q_j e^- \bar{\nu}_e$ occurring inside the baryon core, the flavour changing vector and axial vector currents can be written as

$$\begin{aligned} \mathbf{V}_{ij}(\mathbf{r}) &= \bar{\psi}_{q_i}(\mathbf{r})\boldsymbol{\gamma}\psi_{q_j}(\mathbf{r}) \\ A^{\mu}_{ij}(\mathbf{r}) &= \bar{\psi}_{q_i}(\mathbf{r})\gamma^5\gamma^{\mu}\psi_{q_j}(\mathbf{r}). \end{aligned} \tag{9}$$

Then following the approaches identical to that of Donoghue and Holstein (1982), Barik *et al* (1985) and Rath and Jena (1989), the weak electric and magnetic form factors g_2 and f_2 can be found to be related to certain integrals involving the quark wave functions in the following way:

$$(g_2/g_1) = \mu_p(g_1^{\mathrm{SU(6)}}/g_1)\left(\frac{M_1+M_2}{2M_p}\right)I^{(-)}_{ij} - \frac{(M_1^2-M_2^2)}{4M_1M_2} \tag{10}$$

$$(f_2/f_1) = \mu_p(g_1^{\mathrm{SU(6)}}/g_1)(g_1/f_1)\left(\frac{M_1+M_2}{2M_p}\right)I^{(+)}_{ij} - \frac{(M_1+M_2)^2}{4M_1M_2} \tag{11}$$

where M_p and μ_p are the mass and the magnetic moment of the proton respectively. The wave function integrals $I^{(\pm)}_{ij}$ are

$$I^{(\pm)}_{ij} = \frac{N_{q_i}N_{q_j}}{I_0}\int_0^{\infty} \mathrm{d}r\, r[f_{q_i}(r)g_{q_j}(r) \pm f_{q_j}(r)g_{q_i}(r)] \tag{12}$$

where

$$I_0 = 2N_u^2\int_0^{\infty}\mathrm{d}r\, r f_u(r)g_u(r) = \frac{-3N_u^2}{\lambda_u}. \tag{13}$$

The axial vector form factor g_1 can also be evaluated to give

$$g_1 = g_1^{SU(6)} I_{ij}^A \tag{14}$$

where

$$I_{ij}^A = N_{q_i} N_{q_j} \int_0^\infty \mathrm{d}r [g_{q_i}(r) g_{q_j}(r) - \tfrac{1}{3} f_{q_i}(r) f_{q_j}(r)]. \tag{15}$$

Equation (14) enables one to replace $(g_1^{SU(6)}/g_1)$ appearing in (10) and (11) by $(I_{ij}^A)^{-1}$ so that one can obtain the weak form factor ratios in terms of the wavefunction integrals in ratio form as $(I_{ij}^{(\pm)}/I_{ij}^A)$. This may in fact minimize any uncertainty in the calculation due to any approximation involving the solutions for the quark wavefunctions. The SU(3)-symmetry breaking effects can in fact be introduced through these wavefunction integrals.

3. Evaluation of the wave-function integrals

Now it is clear that in a definite realistic quark model of baryons, one can calculate SU(3)-breaking effects coming from the quark-wave function integrals in the form of $I_{ij}^{(\pm)}$ in (12) and thereby evaluate the weak magnetic and electric form factors in terms of the ratio factors. It can be further noted that by evaluating the wave function integral I_{ij}^A in (15), the axial form factor g_1 can be evaluated in order to obtain independently the values of the weak electric form factor g_2. In the present scheme we retain the SU(2) symmetry in the (u, d) quark sector and break the SU(3)-symmetry in the strange quark (s)-sector only so as to have $m_u = m_d \neq m_s$. Therefore, for the strangeness conserving ($\Delta S = 0$) baryon decays like the neutron beta decay, the wave function integrals would have their trivial values such as

$$I_{uu}^A = \tfrac{1}{3}(4N_u^2 - 1), \quad I_{uu}^{(+)} = 1, \quad I_{uu}^{(-)} = 0. \tag{16}$$

However, for $\Delta S = 1$ transitions they would have different values which are needed to be evaluated in the framework of the present potential model as discussed below.

According to our earlier work (Jena and Panda 1990a) if one considers the average confining potential $V_q(r)$ of the independent constituent quarks in a baryon to be of the form given by (1), then the upper component $g_q(r)$ in (8) satisfies the equation

$$g_q''(r) + \lambda_q[(E_q - m_q - V_0) - a^2 r] g_q(r) = 0 \tag{17}$$

whereas the lower component $f_q(r)$ can always be expressed as

$$f_q(r) = \frac{1}{\lambda_q} r \frac{\mathrm{d}}{\mathrm{d}r}[g_q(r)/r]. \tag{18}$$

Here $\lambda_q = E_q + m_q$. Now choosing a suitable length scale

$$r_{oq} = (\lambda_q a^2)^{-1/3}. \tag{19}$$

Equation (17) can be expressed in terms of a dimensionless variable $\rho = r/r_{oq}$ as

$$g_q''(\rho) + (\varepsilon - \rho) g_q(\rho) = 0 \tag{20}$$

where

$$\varepsilon=(\lambda_q/a^4)^{1/3}(E_q-m_q-V_0). \tag{21}$$

The over-all normalization constant N_q of $\psi_q(\mathbf{r})$ can be easily obtained as

$$N_q=[1+(E_q-m_q-V_0-a^2\langle\langle r\rangle\rangle_q)/\lambda_q]^{-1/2} \tag{22}$$

where $\langle\langle r\rangle\rangle_q$ stands for the expectation value with respect to the normalized radial angular part of the upper component of $\psi_q(\mathbf{r})$. Equation (20) can be solved by reducing it to Airy equation as has been done in our earlier work (Jena and Panda 1990a). Since the Airy function solution for $g_q(\rho)$ makes the evaluation of wavefunction integrals complicated, we use for simplicity the WKB solution which yields

$$g_q(\rho)=\frac{A_q}{(\varepsilon-\rho)^{1/4}}\cos\phi_q(\rho) \tag{23}$$

where

$$\phi_q(\rho)=\int_0^{\rho}\mathrm{d}\rho'(\varepsilon-\rho')^{-1/2}-\pi/4 \tag{24a}$$

$$\varepsilon=(9\pi/8)^{2/3} \tag{24b}$$

and the normalization constant A_q is given by

$$A_q=(r_{oq}^2\varepsilon)^{-1/4}. \tag{24c}$$

Now with the WKB solution available for the upper component $g_q(r)$, one can evaluate the wavefunction integrals $I_{ij}^{(\pm)}$ and I_{ij}^{A} given by (12) and (15) respectively.

If one uses (18) and (20) in (12) then integration by parts can yield

$$I_{us}^{(\pm)}=\pm\frac{N_s}{N_u}[C^{(\pm)}j_0/2-C^{(\mp)}(j_s-j_u)/12] \tag{25}$$

where

$$C^{(\pm)}=(\lambda_u/\lambda_s\pm 1) \tag{26}$$

$$j_0=\int_0^{\infty}\mathrm{d}r\,g_u(r)g_s(r) \tag{27}$$

and

$$j_{q_i}=b_{q_i}^2\int_0^{\infty}\mathrm{d}r\,r^2(\varepsilon-b_{q_i}r)g_u(r)g_s(r) \tag{28}$$

where $b_{q_i}=1/r_{oq_i}$. Again using (18) and (20) in (15) and integrating by parts one can obtain

$$I_{us}^{A}=N_uN_s(j_0-j_1/3\lambda_u\lambda_s r_{os}^2) \tag{29}$$

where

$$j_1=\int_0^{\infty}\mathrm{d}r(\varepsilon-b_s r)g_u(r)g_s(r). \tag{30}$$

Now using WKB solution (23) for $g_q(r)$ and reasonable approximation $\cos\phi_{q_i}\cos\phi_{q_j}\simeq$

$\cos^2\phi_q = \frac{1}{2}$, one can get the integrals (27), (28) and (30) as

$$j_0 = \tfrac{1}{2}KB(1,\tfrac{3}{4})F(\tfrac{1}{4},1,\tfrac{7}{4},K^2) \tag{31}$$

$$j_u = \tfrac{1}{2}K^5\varepsilon^3 B(3,\tfrac{3}{4})F(-\tfrac{3}{4},3,\tfrac{15}{4},K^2) \tag{32}$$

$$j_s = \tfrac{1}{2}K\varepsilon^3 B(3,\tfrac{7}{4})F(\tfrac{1}{4},3,\tfrac{19}{4},K^2) \tag{33}$$

$$j_1 = \tfrac{1}{2}K\varepsilon B(1,\tfrac{7}{4})F(\tfrac{1}{4},1,\tfrac{11}{4},K^2) \tag{34}$$

where $K^2 = r_{os}/r_{ou}$, $B(m,n)$ is the beta function and $F(\alpha,\beta,\gamma,z)$ is the hypergeometric function. Knowledge of the potential parameters and the quark masses together with the ground-state energy eigenvalue E_q can enable one to evaluate the integrals in (31)–(34) which in turn would yield the wave-function integrals $I_{us}^{(\pm)}$ and $I_{us}^{(A)}$ and hence g_2 and f_2.

4. Results and conclusion

The evaluation of form factor ratios (g_2/g_1) and (f_2/f_1) depend on the computation of the integrals $I_{ij}^{(\pm)}$ and I_{ij}^A which on the other hand depend on the potential and quark mass parameters (a, V_0, m_q) and also on certain quantities obtained from the ground state energy eigenvalue solutions for the independent quarks in baryons. These quantities together with the calculated values of the wave-function integrals $I_{ij}^{(\pm)}$ and I_{ij}^A are presented in table 1. The choice of the model parameters are in fact made while using the present model to much wider sector to study the nucleon form factors together with the static properties of the octet baryons in a separate work communicated elsewhere (Jena and Panda 1990b).

Now using the appropriate values of $(I_{ij}^A)^{-1}$, $I_{ij}^{(-)}$ and μ_p from table 1, and the physical baryon masses involved, the form factor ratio (g_2/g_1) for various semileptonic

Table 1. Model parameters and certain quantities corresponding to the ground-state solutions of independent quarks together with the calculated values of the wave-function integrals $I_{ij}^A, I_{us}^{(+)}, I_{us}^{(-)}$.

	Present model $\frac{1}{2}(1+\gamma^0)(a^2r + V_0), a > 0$
Potential parameters	$(a, V_0) = (343, -506)$ MeV
$(m_u = m_d, m_s)$	(243·8, 446·2) MeV
(λ_u, λ_s)	(634·8, 956·2) MeV
(N_u, N_s)	(0·863, 0·913)
(r_{ou}, r_{os})	(2·0715, 1·8071)
(A_u, A_s)	(0·525, 0·562)
K^2	0·8724
ε	2·33811
$(I_{us}^{(+)}, I_{us}^{(-)})$	(0·6812, 0·1134)
(I_{uu}^A, I_{us}^A)	(0·66, 0·544)
(μ_p, μ_n)	$(2·793, -1·913)\mu_N$

Table 2. Calculated values of the weak-electric form factor ratio (g_2/g_1) in comparison with the corresponding predictions in non-relativistic (NR) quark model and MIT bag model.

Decay modes	(g_2/g_1) Present work	(g_2/g_1) (Donoghue and Holstein 1982) NR	MIT
$n\rightarrow pe^-\bar{\nu}_e$	$-0{\cdot}0007$	~ 0	~ 0
$\Lambda\rightarrow pe^-\bar{\nu}_e$	0·5503	0·72	0·291
$\Sigma^-\rightarrow ne^-\bar{\nu}_e$	0·5406	0·72	0·275
$\Xi^-\rightarrow \Lambda e^-\bar{\nu}_e$	0·6711	0·88	0·371
$\Xi^-\rightarrow \Sigma^0 e^-\bar{\nu}_e$	0·7286	0·94	0·415
$\Xi^-\rightarrow \Sigma^+ e^-\bar{\nu}_e$	0·7263	0·94	0·415

decay modes are calculated from (10) and the results are presented in table 2. The corresponding values for (g_2/g_1) obtained by Donoghue and Holstein (1982) in the non-relativistic quark model and MIT bag model are also provided in table 2 for comparison. We observe that our model predictions for (g_2/g_1) are not very much different from the non-relativistic predictions of Donoghue and Holstein but are different from their MIT bag model predictions. Our model estimates for (g_2/g_1) also compare well with the corresponding values obtained in a relativistic power-law potential model by Barik *et al* (1985) and in a relativistic logarithmic potential model by Rath and Jena (1989). Nevertheless it is clear from all these quark model predictions that the sign of the nonvanishing 'g_2' is positive unlike what was envisaged in the earlier studies of Garcia (1971) and Calaprice *et al* (1975). We further notice that consistent with the usual anticipation, the (g_2/g_1) value for $\Delta S = 0$ decay modes such as neutron beta decay is negligibly small while it is appreciable for $\Delta S = 1$ decay modes.

Next, using the appropriate values of $(I_{ij}^A)^{-1}$, $I_{ij}^{(+)}$ and the masses M_1 and M_2 in (11) the form factor ratios (f_2/f_1) for various decay modes are computed in terms of the ratio (g_1/f_1) and are displayed, in table 3. The (g_1/f_1) values are determined in our model from the following considerations. The SU(3) symmetry breaking is usually expected to be a second order effect in (f_2/f_1). Therefore, assuming at the outset that the departures of (f_2/f_1) values from the corresponding Cabibbo predictions are negligibly small, we extract the (g_1/f_1) values from our model expressions provided in table 3. These values when compared with the experimental values of (g_1/f_1) (Bourguin *et al* 1982, 1983) obtained from a Cabibbo fit give quite reasonable agreement (table 3). However, such extracted values of (g_1/f_1) in the Cabibbo limit would correspond to $g_2 = 0$. But it has been observed (Bourguin *et al* 1983) that in order to take into account nonzero values of g_2 the (g_1/f_1) values in the Cabibbo limit are to be modified according to

$$(g_1/f_1)_{g_2\neq 0} = (g_1/f_1)_{g_2=0}[1 - 0{\cdot}24(g_2/g_1)]^{-1}. \tag{35}$$

Incorporating our (g_2/g_1) results from table 2 in (35) we calculate the modified values of $(g_1/f_1)_{g_2\neq 0}$ which are then used in (11) to obtain the weak magnetic form factor ratios (f_2/f_1) in our model. These results are presented in table 4 which compare well with the corresponding values calculated following the work of Donoghue and Holstein (1982). We further notice that the departure of (f_2/f_1) values from the

Table 3. Weak magnetic form factor ratio (f_2/f_1) in terms of (g_1/f_1) and the extracted values of (g_1/f_1) in the Cabibbo limit in comparison with the experimental values ($\mu_p = 2{\cdot}793$ nm, $\mu_n = -1{\cdot}913$ nm).

Decay mode	(f_1/f_2) Present work	(f_2/f_1) Cabibbo	Extracted values of (g_1/f_1) in the Cabibbo limit	(g_1/f_1) Expt. (Bourguin *et al* 1982, 1983)
$n \to pe^- \bar{\nu}_e$	$4{\cdot}230(g_1/f_1) - 1$	$\mu_p - \mu_n - 1$	1·113	1·24
$\Lambda \to pe^- \bar{\nu}_e$	$3{\cdot}827(g_1/f_1) - 1{\cdot}0075$	$\mu_p - 1$	0·732	$0{\cdot}70 \pm 0{\cdot}03$
$\Sigma^- \to ne^- \bar{\nu}_e$	$3{\cdot}982(g_1/f_1) - 1{\cdot}0148$	$\mu_p + 2\mu_n - 1$	−0·258	$-0{\cdot}34 \pm 0{\cdot}05$
$\Xi^- \to \Lambda e^- \bar{\nu}_e$	$4{\cdot}541(g_1/f_1) - 1{\cdot}0072$	$\mu_p + \mu_n - 1$	0·195	$0{\cdot}25 \pm 0{\cdot}05$
$\Xi^- \to \Sigma^0 e^- \bar{\nu}_e$	$4{\cdot}684(g_1/f_1) - 1{\cdot}0026$	$\mu_p - \mu_n - 1$	1·005	1·24
$\Xi^- \to \Sigma^+ e^- \bar{\nu}_e$	$4{\cdot}678(g_1/f_1) - 1{\cdot}0028$	$\mu_p - \mu_n - 1$	1·007	1·24

Table 4. $(g_1/f_1)_{g_2 \neq 0}$ and calculated values of (f_2/f_1) in the present model in comparison with the results of Donoghue and Holstein (1982).

	Present work		Donoghue and Holstein 1982		
				(f_2/f_1)	
Decay mode	$(g_1/f_1)_{g_2 \neq 0}$	(f_2/f_1)	$(g_1/f_1)_{g_2 \neq 0}$	NR	MIT
$n \to pe^- \bar{\nu}_e$	1·113	3·708	1·24	3·625	3·625
$\Lambda \to pe^- \bar{\nu}_e$	0·843	2·219	0·835	1·695	1·9
$\Sigma^- \to ne^- \bar{\nu}_e$	−0·296	−2·193	0·28	−1·97	−2·004
$\Xi^- \to \Lambda e^- \bar{\nu}_e$	0·232	0·046	0·28	0·0652	0·146
$\Xi^- \to \Sigma^0 e^- \bar{\nu}_e$	1·218	4·703	1·392	4·51	4·926
$\Xi^- \to \Sigma^+ e^- \bar{\nu}_e$	1·220	4·704	1·392	4·51	4·926

Cabibbo predictions generated in this manner due to the SU(3) breaking effects does not seem to be sufficiently large so as to be noticeable in a conventional Cabibbo fit.

The over all agreement of the results of our model with some of those due to Donoghue and Holstein (1982), Barik *et al* (1985), and Rath and Jena (1989) gives only a desirable feature to justify its relevance in the study of hadronic properties. In fact, there exists in the literature, various types of phenomenological potential models, which more or less agree with each other in describing different hadronic properties like the mass spectra, the static electromagnetic properties, and nucleon form factors, etc. Although each of these models has varying degrees of theoretical and phenomenological justifications in its favour, the existing data does not seem to prefer uniquely any one of them. Only a wider range of data involving certain aspects more sensitive to the specific features of the models may enable one to select out some in preference over others. In that respect, the weak electric form factor ratio (g_2/g_1) for various semileptonic decays of the baryons may provide a guiding factor.

Acknowledgements

We are thankful to Prof. N Barik of the Department of Physics, Utkal University for useful discussions. We also gratefully acknowledge the valuable suggestions of Prof. B B Deo of Berhampur University.

References

Barik N, Dash B K and Das M 1985 *Phys. Rev.* **D32** 1725
Bourguin M *et al* 1982 *Z. Phys.* **C12** 307
Bourguin M *et al* 1983 *Z. Phys.* **C21** 1, 17, 27
Bohm A, Kielanowski P and Garcia A 1984 *Phys. Rev.* **D30** 231
Brandle H *et al* 1978 *Phys. Rev. Lett.* **40** 306
Cabibbo N 1963 *Phys. Rev. Lett.* **10** 531
Calaprice F P, Freedmann S J, Mead W C and Vantite H C 1975 *Phys. Rev. Lett.* **35** 1566
Donoghue J F and Johnson K 1980 *Phys. Rev.* **D21** 1975
Donoghue J F and Holstein B R 1982 *Phys. Rev.* **D25** 206
Eich E, Rein D and Rodenberg R 1985 *Z. Phys.* **C28** 225
Garcia A 1971 *Phys. Rev.* **D3** 2638
Garcia A and Kielanowski P 1982 *Phys. Lett.* **B110** 498
Jena S N and Panda S 1990a *Pramana – J. Phys.* **35** 21
Jena S N and Panda S 1990b Berhampur University – Preprint
Keppinger W, Calaprice F and Miller D 1978 *Bull. Am. Phys. Soc.* **23** 603
Kobudera K *et al* 1985 *Nucl. Phys.* **A439** 695
Lebrun P *et al* 1978 *Phys. Rev. Lett.* **40** 302
Rath D P and Jena S N 1989 *Pramana – J. Phys.* **32** 753
Sugimoto K, Tanihita I and Goring J 1975 *Phys. Rev. Lett.* **34** 1533
Tribble R E and May D P 1978 *Phys. Rev.* **C18** 2704