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Hypernuclear currents in a relativistic mean-field theory with tensor ωYY couplings

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The possibility of strong tensor coupling of the (isoscalar) $\omega(783)$ meson to hyperons, specifically $\Lambda(1115)$, has been suggested by several authors during the past decade. This paper studies the effect of this coupling on magnetic properties of hypernuclei, within the framework of the (minimal) $\sigma+\omega$ model. Our results are analytical and intuitive. They provide physical insight into recent numerical results of Gattone, Chiapparini, and Izquierdo. Additionally, we point out similar effects in transition form factors at moderate momentum transfer, on finite nuclei.

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It has been pointed out some time ago [1] that a self-consistent calculation of hypernuclear currents yields an interesting difference between the nonrelativistic and a relativistic σ-ω model predictions for observables such as Λ-hypernuclear magnetic moments. The nonrelativistic model considered was an extreme single-particle theory where the magnetic moments are simply the Schmidt values; the relativistic model takes into account both the hyperon single-particle (baryon) current and the linear response of the core due to the hyperon (in Ref. [1] this latter effect has been calculated in the local density approximation).

In Ref. [1] only pure scalar and vector meson-baryon couplings have been considered. However it has since been pointed out [2,3] that there is a strong theoretical support for an additional ωYY coupling, of the tensor type. While the ωNN tensor vertex is consistently negligible, thus leaving intact the success of the Dirac many-body theory for regular nuclei, the ωYY (and in particular the $\omega\Lambda\Lambda$) tensor coupling is large. It is therefore reasonable to start with a mean-field theory (MFT) Lagrangian:

$$\begin{aligned} L_{\rm MFT} = {} & \bar\psi_N[\gamma_\mu(i\partial^\mu - g_v^N V^\mu) - (M_N - g_s^N\phi_0)]\psi_N \\ & + \bar\psi_Y[\gamma_\mu(i\partial^\mu - g_v^Y V^\mu) - (M_Y - g_s^Y\phi_0)]\psi_Y \\ & - \frac{f_{\omega YY}}{4M_Y}\bar\psi_Y\sigma^{\mu\nu}F_{\mu\nu}\psi_Y + \tfrac{1}{2}(\partial_\mu\phi_0\partial^\mu\phi_0 - m_s^2\phi_0^2) \\ & - \tfrac{1}{4}F_{\mu\nu}F^{\mu\nu} + \tfrac{1}{2}m_v^2 V_\mu V^\mu \ . \end{aligned} \quad (1)$$

In Eq. (1) ψ_N and ψ_Y represent the Dirac spinors for the nucleons and Y hyperon, ϕ_0 is the scalar field, and $V^\mu=(V^0,\mathbf{V})$ is the vector field (note that maintaining the three-vector component of V^μ is crucial in our discussion), $f_{\omega YY}$ is the tensor ωYY coupling constants as described in Ref. [3] (and we use here the result $f_{\omega NN} \ll f_{\omega YY}$), and

$$F_{\mu\nu} \equiv \partial_\mu V_\nu - \partial_\nu V_\mu \ . \quad (2)$$

In Ref. [3] the tensor term was derived from a quark model. In a purely hadronic theory such terms are sometimes put by hand into the Lagrangian (in analogy with magnetic-moment contributions of Pauli type to represent the interaction of the electromagnetic field with the anomalous magnetic moment of the baryon). It has long been realized that such contributions should emerge (in terms of hadronic degrees of freedom) as higher-order perturbative corrections in the field-theoretical calculations [4]. However, a complete theory along these lines has never been worked out. It is important, though, to keep in mind that the magnetic moments and tensor couplings should not be treated as tree-level contributions in the strict hadronic theory sense.

Lagrange's equations (the equations of motion) derived from the Lagrangian, Eq. (1), differ from standard [5] MFT equations by the addition of a tensor coupling for the $\omega\Lambda\Lambda$ vertex; thus, the hyperon equation of motion is

$$\left[\gamma_\mu(i\partial^\mu - g_v^N V^\mu) - (M_Y - g_s^Y\phi_0) - f_{\omega YY}\frac{1}{4M_Y}\sigma^{\mu\nu}F_{\mu\nu}\right]\psi_Y = 0 \ , \quad (3)$$

while that of the ω is

$$\partial_\mu F^{\mu\nu} + m_v^2 V^\nu = g_v^N\bar\psi_N\gamma^\nu\psi_N + g_v^Y\bar\psi_Y\gamma^\nu\psi_Y - \frac{f_{\omega YY}}{2M_Y}\frac{\partial}{\partial x^\mu}(\bar\psi_Y\sigma^{\mu\nu}\psi_Y) \ . \quad (4)$$

In Eq. (4) we identify a two-part conserved baryon current:

$$\begin{aligned} j_B^\nu & \equiv j_{B_N}^\nu + j_{B_Y}^\nu \\ & \propto g_N^v\bar\psi_N\gamma^\nu\psi_N \\ & \quad + \left[g_v^Y\bar\psi_Y\gamma^\nu\psi_Y - \frac{f_{\omega YY}}{2M_Y}\frac{\partial}{\partial x^\mu}(\bar\psi_Y\sigma^{\mu\nu}\psi_Y)\right] \ . \end{aligned} \quad (5)$$

Note the presence of the ωYY tensor coupling in Eqs. (1) and (3)–(5). Gattone, Chiapparini, and Izquierdo [6] have recently calculated numerically the effect of this term on the results of Refs. [1] and [7]. The purpose of this paper is to provide a simple explanation of the results obtained in [6] and throw more light on the physics involved in it.

Toward this end we start with the *nucleon* Dirac equation,

$$[-i\boldsymbol{\alpha}\cdot\nabla + \beta(M_N - g_s^N\phi_0) + g_v^N V_0 - \boldsymbol{\alpha}\cdot g_v^N\mathbf{V}]\psi_N = E_N\psi_N \quad (6)$$

which gives after a standard manipulation [8]

$$2(E_N-g_v^N V_0)\psi_N^\dagger(\mathbf{r})\boldsymbol{\alpha}\psi_N(\mathbf{r}) = i\psi_N^\dagger(\mathbf{r})(\overleftarrow{\nabla}-\overrightarrow{\nabla})\psi_N(\mathbf{r}) + \overrightarrow{\nabla}\times[\psi_N^\dagger(\mathbf{r})\boldsymbol{\Sigma}\psi_N(\mathbf{r})]-2g_v^N\mathbf{V}\psi_N^\dagger\psi_N \ . \quad (7)$$

Here, $\boldsymbol{\Sigma}=(\begin{smallmatrix}\sigma & 0\\ 0 & \sigma\end{smallmatrix})$. The first two terms on the right-hand side of Eq. (7) are convection and spin contributions, while the third one results from the spacelike vector potential **V**. Note that in the nucleon case there is no other spin contribution as $f_{\omega NN}\simeq 0$.

The spinors ψ_N in Eqs. (6) and (7) are the solutions for the bound state single-particle wave functions for the entire system (nuclear core+hyperon), not just the closed-core (spherical) problem. Equation (7) can be compared with the Gordon decomposition for free, on-shell nucleons,

$$\gamma_\mu \to \frac{(p+p')_\mu}{2M_N}+\frac{i}{2M_N}\sigma_{\mu\nu}q^\nu \ , \quad (8)$$

where p,p' are the initial and final N four-momenta and q is the momentum transfer $q=p'-p$. Equation (8) is used by Gattone, Chiapparini, and Izquierdo [6] in an approximate finite nucleus calculation. It is also intuitively useful to recall that in infinite nuclear matter the spinors satisfy [5]

$$u^\dagger u=\frac{E_\kappa^*}{M^*}\bar{u}u \quad (9)$$

with $M^*=M-g_s\phi_0$, $E_\kappa^*=(\kappa^2+M^{*2})^{1/2}$, $\boldsymbol{\kappa}=\mathbf{k}-g_v\mathbf{V}$. Equation (7) can then be described as the Gordon decomposition with scalar and vector interactions included. It explicitly shows the spin, orbital, and vector potential contributions.

In the presence of interactions, M is replaced by M^*; this would seem to create the usual M/M^* enhancement, however, the core response effect (where the presence of **V** is crucial) will, in general, cancel it. In Ref. [1], it has been shown that this cancellation is not complete when dealing with Λ hypernuclei; however, Gattone, Chiapparini, and Izquierdo [6] subsequently showed numerically that when the $\omega\Lambda\Lambda$ tensor coupling [2,3] is included, an almost complete cancellation along these lines is again achieved for the ground state (Λ in the 1s state).

We present here an intuitive, physically transparent derivation of the numerical results of Ref. [6]. Building upon published work, we feel that it is unnecessary to provide here all technical details, although we will try to provide a self-contained discussion. The three-vector **V** in Eqs. (6) and (7) is related to the baryon current through

$$\mathbf{V}=\frac{g_v^N}{m_v^2}\langle \mathbf{j}_N\rangle+\frac{g_v^Y}{m_v^2}\langle \mathbf{j}_Y\rangle \ , \quad (10)$$

where the individual components of the current, $\langle \mathbf{j}_N\rangle$ and $\langle \mathbf{j}_Y\rangle$, are proportional to the pure vector parts of the currents identified in Eq. (5); using Eq. (7) $\langle \mathbf{j}_N\rangle$ can be expressed as

$$\langle \mathbf{j}_N\rangle=\sum_N \psi_N^\dagger\boldsymbol{\alpha}\psi_N = \sum_N\frac{1}{2(E_N-g_v^N V_0)}[i\psi_N^\dagger(\overleftarrow{\nabla}-\overrightarrow{\nabla})\psi_N+\overrightarrow{\nabla}\times(\psi_N^\dagger\boldsymbol{\Sigma}\psi_N)] - \sum_N\frac{g_v^N\mathbf{V}}{E_N-g_v^N V_0}\psi_N^\dagger\psi_N \ . \quad (11)$$

We now assume that the energy denominators can be taken outside of the sums, using an average energy $\bar{E}_N$. The first term vanishes for closed-shell nucleon configurations (with total orbital and spin angular momenta equal to 0). Moreover, in calculating the nuclear magnetic moment from the electromagnetic currents $\mathbf{j}^{\rm em}$

$$\bar{\mu}\propto\int d\mathbf{r}\tfrac{1}{2}\mathbf{r}\times(\mathbf{j}_N^{\rm em}+\mathbf{j}_Y^{\rm em}) \ ,$$

the orbital contribution would vanish for closed shell, $L=0$ configurations. Although Eq. (10) is exact only for nuclear matter, we *assume* its validity in this work.

The second, nonvanishing term in Eq. (11) gives

$$\langle \mathbf{j}_N\rangle=-g_v^N\mathbf{V}\frac{\langle\rho_N\rangle}{\bar{E}_N-g_v^N V_0} \ , \quad (12)$$

where ρ_N is the vector density of nucleons (the quantity usually referred to as "density" in "classical" nuclear physics). Note that $\langle \mathbf{j}_N\rangle$ does not vanish since the spinors ψ_N in Eq. (11) are solutions of the one-body equation (6) for the whole-nucleus model (and not just the spherically symmetric closed core).

We can now use Eq. (10) to solve for $\langle \mathbf{j}_N\rangle$:

$$\langle \mathbf{j}_N\rangle=-\left[1+\left[\frac{g_v^N}{m_v}\right]^2\frac{\langle\rho_N\rangle}{\bar{E}_N-g_v^N V_0}\right]^{-1}\frac{g_v^N g_v^Y}{m_v^2}\frac{\langle\rho_N\rangle}{\bar{E}_N-g_v^N V_0}\langle \mathbf{j}_Y\rangle \ . \quad (13)$$

[Note that the term omitted in going from Eq. (11) to Eq. (12) due to its vanishing contribution could be included in Eq. (13) for non-closed-shell situations, however in the rest of this paper its contribution is assumed to vanish.]

The result in Eq. (13) is similar to the one obtained in Ref. [1], since the total baryon current is $\langle \mathbf{j}_B\rangle=\langle \mathbf{j}_N\rangle+\langle \mathbf{j}_Y\rangle$, or

$$\langle \mathbf{j}_B\rangle=\left[1+\left[\frac{g_v^N}{m_v}\right]^2\frac{\langle\rho_N\rangle}{\bar{E}_N-g_v^N V_0}\right]^{-1}\left[1+\frac{g_v^N}{m_v^2}(g_v^N-g_v^Y)\frac{\langle\rho_N\rangle}{\bar{E}_N-g_v^N V_0}\right]\langle \mathbf{j}_Y\rangle \ . \quad (14)$$

This result gave rise to a relativistic effect when compared with the nonrelativistic extreme single-particle Schmidt theory. The advantage of the present derivation over that of Ref. [1] is that it allows us to take into account the $\omega\Lambda\Lambda$ tensor coupling (which is not possible in a model of infinite nuclear matter).

As already introduced in Eq. (5), the ωYY tensor coupling results in a modification of the baryon current whereby the vector current

$$\psi_Y^\dagger \boldsymbol{\alpha}\psi_Y=\frac{1}{2(E_Y-g_v^Y V_0)}[i\psi_Y^\dagger(\overleftarrow{\nabla}-\vec{\nabla})\psi_Y+\vec{\nabla}\times(\psi_Y^\dagger \boldsymbol{\Sigma}\psi_Y)] \tag{15}$$

is modified by adding to it the tensor term

$$\frac{f_{\omega YY}}{g_v^Y}\frac{1}{2M_Y}\vec{\nabla}\times(\psi_Y^\dagger\beta\boldsymbol{\Sigma}\psi_Y) . \tag{16}$$

The extra (tensor) part is obviously expected to modify the results of Refs. [1,7]. The pertinent physical observable dealt with in Refs. [1,6,7] is the Λ-hypernuclear magnetic moment.

The nuclear magnetic moment has both core and Λ (the single hyperon outside of the closed nucleon core) contributions. For clarity of presentation we deal with an s-state Λ (no orbital contributions) first, and then with $l>0$ Λ single-particle states. For s states ($l=0$), which are described in terms of real wave functions, the first term on the right-hand side of Eq. (15), namely, the one containing the $\overleftarrow{\nabla}-\vec{\nabla}$ operator, does not contribute to the Λ baryon current.

In calculating the hypernuclear magnetic moment,

$$\boldsymbol{\mu}=\sum^{A}\int d^3\mathbf{r}\tfrac{1}{2}\mathbf{r}\times\mathbf{j}^{\rm em}(\mathbf{r}) \tag{17}$$

it is straightforward to write down the hypernuclear electromagnetic current $\mathbf{j}^{\rm em}$ based on the preceding discussion. Since the Λ is neutral, it contributes to $\mathbf{j}^{\rm em}$ only through its anomalous magnetic moment μ_Λ (no orbital contribution),

$$\mathbf{j}_\Lambda^{\rm em}(\mathbf{r})=\frac{\mu_\Lambda}{2M_N}\vec{\nabla}\times[u_\Lambda^\dagger(\mathbf{r})\beta\boldsymbol{\Sigma}u_\Lambda(\mathbf{r})] . \tag{18}$$

This contribution is very close to the Schmidt value with deviations of the order of $O(F_\Lambda^2/G_\Lambda^2)$, where F_Λ and G_Λ are the lower and upper components of the Λ Dirac spinor.

The core (closed nucleon shells) contribution,

$$\mathbf{j}_{\rm core}^{\rm em}=\tfrac{1}{2}\sum_N\psi_N^\dagger\boldsymbol{\alpha}\psi_N , \tag{19}$$

is identical in form to the nucleon baryon current in Eqs. (11)–(16). The nuclear magnetic moment for a many-body state with total angular moment J is defined as

$$\mu\equiv\langle\mu_Z\rangle_{\rm max}-\frac{\langle\boldsymbol{\mu}\cdot\mathbf{J}\rangle}{\langle J^2\rangle}J .$$

In Refs. [1,7] deviations from the Schmidt values were obtained when only a pure vector $\omega\Lambda\Lambda$ coupling was used in the core current, Eq. (19). In this work we are dealing (following Gattone *et al.* [6]) with $\omega\Lambda\Lambda$ tensor coupling, Eq. (16), as well. Thus, the core contribution to the magnetic moment Eq. (17) is obtained from Eqs. (13) and (19) as

$$\boldsymbol{\mu}_{\rm core}=-\frac{1}{4}\frac{g_v^N g_v^\Lambda}{m_v^2}\int d^3r\left[1+\left[\frac{g_v^N}{m_v}\right]^2\frac{\rho_N(r)}{\bar{E}_N-g_v^N V_0(r)}\right]^{-1}\frac{\rho_N(r)}{\bar{E}_N-g_v^N V_0(r)}\mathbf{r}$$
$$\times\left\{\frac{1}{2[E_\Lambda-g_v^\Lambda V_0(r)]}\{i\psi_\Lambda^\dagger(\mathbf{r})(\overleftarrow{\nabla}-\vec{\nabla})\psi_\Lambda(\mathbf{r})+\vec{\nabla}\times[\psi_\Lambda^\dagger(\mathbf{r})\boldsymbol{\Sigma}\psi_\Lambda(\mathbf{r})]\}+\frac{f_{\omega\Lambda\Lambda}}{g_v^\Lambda}\frac{1}{2M_\Lambda}\vec{\nabla}\times(\psi_\Lambda^\dagger\beta\boldsymbol{\Sigma}\psi_\Lambda)\right\} . \tag{20}$$

In Eq. (20), it is possible to replace the density ρ_N by the vector potential using

$$\rho_N=\frac{m_v^2}{g_v^N}V_0 .$$

For an s-wave Λ the term containing the $\vec{\nabla}-\overleftarrow{\nabla}$ operator vanishes. Furthermore we may approximate as follows, once more using the methods of [8]:

$$\psi_\Lambda^\dagger\beta\boldsymbol{\Sigma}\psi_\Lambda\equiv\frac{M_\Lambda-g_s^\Lambda\phi_0}{E_\Lambda-g_v^\Lambda V^0}\psi_\Lambda^\dagger\boldsymbol{\Sigma}\psi_\Lambda-\frac{1}{2(E_\Lambda-g_v^\Lambda V^0)}[i\nabla(\psi_\Lambda^\dagger\beta\gamma^5\psi_\Lambda)+\psi_\Lambda^\dagger\beta\gamma^5\boldsymbol{\Sigma}\times(\vec{\nabla}-\overleftarrow{\nabla})\psi_\Lambda]+\frac{f_{\omega\Lambda\Lambda}}{M_\Lambda g_v^\Lambda}\psi_\Lambda^\dagger\gamma^5\boldsymbol{\Sigma}\times(\vec{\nabla}g_v^\Lambda V^0)\psi_\Lambda$$
$$\approx\frac{M_\Lambda^*}{E_\Lambda^*}\psi_\Lambda^\dagger\boldsymbol{\Sigma}\psi_\Lambda+O(v^2/c^2) , \tag{21}$$

where v is the typical velocity of the Λ in the nucleus. Since the average Λ kinetic energy is ≈20 MeV, $v^2/c^2\approx0.03$–0.06 (depending on the effective mass M_Λ^*). The three (vector) terms in Eq. (21), dismissed as $O(v^2/c^2)$, could conceivably reinforce, giving a 10–20% correction. Since, however, these vectors generally point in quite different directions and are cross multiplied by $\mathbf{r}$ it seems likely the net correction will be smaller than 5% in practical calculations, comparable to the deviation of M_Λ^*/M_Λ from unity in the nuclear surface (which we neglect below). Neglecting the terms of $O(v^2/c^2)$ in Eq. (21) we find

$$\boldsymbol{\mu}_{\rm core}\simeq-\frac{1}{4}g_v^\Lambda\int d^3r\left[1+g_v^N\frac{V_0(r)}{\bar{E}_N-g_v^N V_0(r)}\right]^{-1}\frac{V_0(r)}{\bar{E}_N-g_v^N V_0(r)}\mathbf{r}\times\left\{\frac{1}{2E_\kappa^{\Lambda*}}\vec{\nabla}\times[\psi_\Lambda^\dagger(\mathbf{r})\boldsymbol{\Sigma}\psi_\Lambda(\mathbf{r})]\right.$$
$$\left.+\frac{f_{\omega\Lambda\Lambda}}{g_v^\Lambda}\frac{1}{2M_\Lambda}\vec{\nabla}\times\frac{M_\Lambda^*}{E_\Lambda^*}[\psi_\Lambda^\dagger(\mathbf{r})\boldsymbol{\Sigma}\psi_\Lambda(\mathbf{r})]\right\} . \tag{22}$$

References [2] and [3] argued that for the Λ hyperon

$$\frac{f_{\omega\Lambda\Lambda}}{g_v^\Lambda}=-1 . \tag{23}$$

[In models that include $q\bar{q}$ pair corrections to the naive quark bag, the tensor coupling need not satisfy Eq. (23), because anomalous terms are expected.] Since the derivatives ($\vec{\nabla}\times$ operator) limit the values of $\mathbf{r}$ to the nuclear surface where difference between M_Λ and M_Λ^* are small, we find from Eq. (22) that

$$\mu(\text{core}) \simeq 0 \ . \tag{24}$$

This interesting result means that for a Λ hyperon in the $1s$ state, and in the presence of $\omega\Lambda\Lambda$ tensor coupling, the core contribution to the hypernuclear magnetic moment is nonexistent or small and the expected result is simply the Schmidt value. This is a correction of previous results [1,7], where the tensor coupling, Eq. (23), has been neglected. The present result, Eq. (24), is in agreement with that of Ref. [6].

For higher orbital angular momentum ($l_\Lambda > 0$) states, orbital contributions arise from the term involving the $\overleftarrow{\nabla} - \vec{\nabla}$ operator [9] in the hyperon current [e.g., Eq. (15)]. These orbital contributions are large and their net contribution to the electromagnetic current and the magnetic moment operator is always negative [1,6]. In order to understand the results of Ref. [6] we note that the last two terms in Eq. (20) (namely, the anomalous—or tensor—type terms) still mutually cancel for $l_\Lambda > 0$. Thus, the nonvanishing contribution comes from the first term, involving the $\overleftarrow{\nabla} - \vec{\nabla}$ operator. However, this term is also the only contributor to a nuclear-matter or a local-density-approximation (LDA) calculation such as Ref. [1] (no magnetic-moment-type contributions exist in nuclear matter).

Thus, for $l_\Lambda > 0$ there is a great similarity between the nuclear matter (or LDA) and the finite-nucleus calculations, when the latter includes the effect of the $\omega\Lambda\Lambda$ tensor coupling. We therefore expect the LDA results of Ref. [1] to be in good agreement with $l_\Lambda > 0$ finite-nucleus calculation of Ref. [6]. The core contribution is negative in both cases and its magnitude is also similar, with differences at a level which can be fully expected when comparing an LDA with a finite-nucleus calculation.

Using similar steps we have also derived the dynamical electromagnetic matrix element between states differing only in the hyperon orbital. The same techniques that led to Eqs. (7), (11), (15), (16), and (20) yield the total electromagnetic current matrix element (for Λ hypernuclei)

$$\begin{aligned}
\langle f|\mathbf{j}^{\text{em}}|i\rangle &= \langle f|\mathbf{j}_\Lambda^{\text{em}}|i\rangle + \langle f|\mathbf{j}_{\text{core}}^{\text{em}}|i\rangle \\
&= \frac{\mu_\Lambda}{2M_N}\vec{\nabla}\times[\psi_{\Lambda_f}^\dagger(\mathbf{r})\beta\boldsymbol{\Sigma}\psi_{\Lambda_i}(\mathbf{r})] \\
&\quad -\frac{1}{2}g_v^\Lambda\frac{V_0(r)}{\bar{E}_N - g_v^N V_0(r)}\left\{\frac{i}{E_{\Lambda_f}+E_{\Lambda_i}-2g_v^\Lambda V_0(r)}\{\psi_{\Lambda_f}^\dagger(\mathbf{r})(\overleftarrow{\nabla}-\vec{\nabla})\psi_{\Lambda_i}(\mathbf{r})+\vec{\nabla}\times[\psi_{\Lambda_f}^\dagger(\mathbf{r})\boldsymbol{\Sigma}\psi_{\Lambda_i}(\mathbf{r})]\}\right. \\
&\qquad \left. +\frac{f_{\omega\Lambda\Lambda}}{g_v^\Lambda}\frac{1}{2M_\Lambda}\vec{\nabla}\times[\bar{\psi}_{\Lambda_f}(\mathbf{r})\boldsymbol{\Sigma}\psi_{\Lambda_i}(\mathbf{r})]\right\} .
\end{aligned} \tag{25}$$

Noting from Eqs. (15), (16), and (20) that the contribution to the Λ *baryon* current ($\mathbf{j}_\Lambda$) is very small, i.e., the last two terms in Eq. (25) cancel each other, we find that our present observations for the hypernuclear magnetic moments also hold for the dynamical electromagnetic current matrix element.

In conclusion, we have demonstrated in an intuitive, analytical manner that the results of Ref. [6] can be understood from those of Ref. [1], upon adding $\omega\Lambda\Lambda$ tensor coupling. We have included finite nucleus, bound-state wave functions for the *full* system (nuclear core plus hyperon) and have used generalizations for finite nuclei of the Gordon decomposition. We constructed and analyzed full nuclear currents (including the induced core contribution), including both static and dynamic matrix elements.

The approximate spin independence of the Λ baryon current depends, of course, on the precise numerical value of Eq. (23). This value would differ for other hyperons [3]. However, magnetic moments are not expected to be measured for the more exotic hypernuclei (Σ, Ξ, etc., hypernuclei) in the forseeable future.

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