Note on spin-orbit interactions in nuclei and hypernuclei

N. Kaiser and W. Weise

Physik Department, Technische Universität München, D-85747 Garching, Germany

Abstract

A detailed comparison is made between the spin-orbit interactions in Λ hypernuclei and ordinary nuclei. We argue that there are three major contributions to the spin-orbit interaction: 1) a short-range component involving scalar and vector mean fields; 2) a "wrong-sign" spin-orbit term generated by the pion exchange tensor force in second order; and 3) a three-body term induced by two-pion exchange with excitation of virtual Δ(1232)-isobars (a la Fujita-Miyazawa). For nucleons in nuclei the long-range pieces related to the pion-exchange dynamics tend to cancel, leaving room dominantly for spin-orbit mechanisms of short-range origin (parametrized e.g. in terms of relativistic scalar and vector mean fields terms). In contrast, the absence of an analogous 2π-exchange three-body contribution for Λ hyperons in hypernuclei leads to an almost complete cancellation between the short-range (relativistic mean-field) component and the "wrong-sign" spin-orbit interaction generated by second order π-exchange with an intermediate Σ hyperon. These different balancing mechanisms between short- and long-range components are able to explain simultaneously the very strong spin-orbit interaction in ordinary nuclei and the remarkably weak spin-orbit splitting in Λ hypernuclei.

PACS: 21.30.Fe, 21.80+a, 24.10.Cn

1 Introduction

The microscopic understanding of the dynamical origin behind the strong nuclear spin-orbit force is still one of the key questions in nuclear physics. The analogy with the spin-orbit interaction in atomic physics gave the hint that it could be a relativistic effect. This idea has lead to the phenomenologically successful scalar-vector mean field models for nuclear structure calculations [1, 2]. In these models the nucleus is described as an ensemble of independent Dirac quasi-particles moving in self-consistently generated scalar and vector mean fields. With nucleon Fermi momenta typically less than a third of the (free) nucleon mass, the motion of the nucleons in nuclei is non-relativistic on average. Signatures of relativity are nonetheless manifest in the large spin-orbit coupling which emerges in that framework naturally from the interplay of the individually large scalar and vector mean fields of opposite signs. Their sum balances such as to produce a relatively weak central potential, whereas their difference coherently generates the strong spin-orbit potential [3, 4, 5].

In the context of QCD sum rule calculations these scalar and vector mean fields can be related to the leading changes of the scalar quark condensate ⟨\bar{q}q⟩ and the quark density ⟨q\dagger q⟩ at finite baryon density. This connection has been utilized in Ref.[6] to derive a relativistic nuclear energy density functional constrained by low-energy QCD. In such an approach nuclear binding originates primarily through pionic fluctuations (i.e. two-pion exchange calculated with in-medium chiral perturbation theory) while the spin-orbit interaction results from the strong scalar

\footnote{Work supported in part by BMBF, GSI and by the DFG cluster of excellence: Origin and Structure of the Universe.}
and vector mean fields related to changes of the condensate structure of the QCD vacuum at finite baryon density. Nuclear structure calculations of spherical and deformed nuclei performed within this approach reach the same level of accuracy, in comparison with data, as the best phenomenological relativistic mean field models.

In a phenomenological boson exchange picture the relativistic spin-orbit interaction is modeled by the exchange of a scalar boson ("$\sigma$") and a vector boson ("$\omega$") between nucleons, simulating short-range dynamics but ignoring the physics at intermediate and long ranges characteristic of the average distance scale between nucleons in nuclei. Of particular interest in the context of the spin-orbit interaction are the effects from the exchange of two pions. As a matter of fact, the dominant tensor interaction from one-pion exchange has a spin- and momentum dependence that produces, in second order, the spin-orbit coupling $i(\vec{\sigma}_1 + \vec{\sigma}_2) \cdot (\vec{q} \times \vec{p})$ in momentum space [7]. Here, $\vec{\sigma}_{1,2}$ are the usual Pauli spin-operators of the nucleons; $\vec{p}$ and $\vec{q}$ denote the center-of-mass momentum and the momentum transfer in the elastic $NN$ scattering process. The explicit calculation of the relevant pion loop integral representing iterated one-pion exchange (the second-order tensor force) gives the following isoscalar spin-orbit $NN$ scattering amplitude at threshold ($\vec{p} = \vec{q} = 0$):

$$V_{so}^{(\pi\pi)} = -\frac{g_A^4 M_N}{64 \pi f^4_\pi m_\pi} < 0,$$

with the nucleon axial vector constant $g_A \simeq 1.3$ and the pion decay constant $f_\pi = 92.4$ MeV. At first sight, this contribution has the "wrong" sign in comparison to the one from vector boson exchange,

$$V_{so}^{(\omega)} = \frac{3g_\omega^2 N}{4M_N^2 m_\omega^2} > 0,$$

or scalar boson exchange,

$$V_{so}^{(\sigma)} = \frac{g_\sigma^2 N}{4M_N^2 m_\sigma^2} > 0.$$

As it stands the "wrong-sign" spin-orbit interaction from iterated $1\pi$-exchange in the two-nucleon system must also show up in nuclear many-body systems.

On the other hand it has long been known that calculations based on Hamiltonians with realistic two-nucleon potentials (fitting accurately all empirical $NN$ phase shifts and mixing angles) cannot correctly predict the observed spin-orbit splittings of nuclear levels in light nuclei [8, 9]. A suitably chosen three-body force was then implemented to account for this deficiency. In fact, one of the basic motivations for the genuine three-nucleon interaction introduced originally by Fujita and Miyazawa [10] was the study of such spin-orbit splittings. The Fujita-Miyazawa three-nucleon interaction involves the exchange of two pions between the three nucleons, where the central nucleon which couples to both pions is excited to a $\Delta(1232)$ resonance in the intermediate state. When considered in the vacuum this $2\pi$-exchange process with excitation of virtual $\Delta(1232)$ isobars generates the long-range tail of the isoscalar central attraction between nucleons. However, in the medium the Pauli blocking of intermediate nucleon states below the Fermi surface changes the (orbital) angular momentum balance and an additional spin-orbit interaction (which increases with density) emerges. This long ranged pion-induced three-body interaction provides an additional important contribution to the spin-orbit force in nuclear many-body systems.
In summarizing this introductory discussion, we thus identify three major sources of the spin-orbit interaction in nuclear systems:

- **Short distance dynamics of coherently acting scalar and vector fields**
  These effects are frequently parametrized in terms of sigma-omega exchange models. Since only the ratios $g^2/m^2$ of squared coupling constants and boson masses appear in any relevant calculated quantity, an equivalent effective field theory approach with $NN$ contact interactions, encoding short distance dynamics with coupling strength parameters $G_V = g^2_{2N}/m^2_\omega$ and $G_S = g^2_{\sigma N}/m^2_\sigma$, will yield identical output. In alternative QCD sum rule descriptions these terms can be thought of as originating from leading in-medium changes of the quark condensate $\langle \bar{q}q \rangle$ and the quark density $\langle q^\dagger q \rangle$.

- **Intermediate and long range spin-orbit dynamics**
  induced by the pion exchange tensor force in second order, with Pauli blocking acting on intermediate nucleon states. This mechanism has an analogous counterpart in hypernuclei where the $\Lambda N$ interaction involves an intermediate $\Sigma$ hyperon between pion exchanges.

- **Three-body spin-orbit interaction**
  of the Fujita-Miyazawa type, generated by two-pion exchange with intermediate excitation of a virtual $\Delta$ isobar. It is this latter important ingredient that has no counterpart in $\Lambda$ hypernuclei. As a consequence, the balance of all three major contributions to the spin-orbit is qualitatively different between nuclei and hypernuclei, as we shall demonstrate.

## 2 Spin-orbit coupling for nucleons

We begin with a reminder of facts about nuclear spin-orbit interactions. In order to quantify the spin-orbit terms discussed before it is convenient to consider the energy density functional which serves as a general starting point for nuclear structure calculations (of heavy nuclei) within the self-consistent mean-field approximation \[11\]. The spin-orbit coupling term in the nuclear energy density functional has the form:

$$ E_{so}[\rho, \vec{J}] = F_{so}(\rho) \vec{\nabla} \rho \cdot \vec{J}. \quad (1) $$

It is constructed from the gradient of the (ordinary) nuclear density distribution $\rho(\vec{r})$ and the so-called spin-orbit density, defined by:

$$ \vec{J}(\vec{r}) = \sum_{\alpha \in \text{occ}} \Psi_\alpha^\dagger(\vec{r}) i \vec{\sigma} \times \vec{\nabla} \Psi_\alpha(\vec{r}). \quad (2) $$

The strength of the nuclear spin-orbit coupling is measured by the density dependent strength function $F_{so}(\rho)$.\footnote{The frequently used single-particle spin-orbit potential strength is $U_{Nls}(\rho) = 2\rho F_{so}(\rho)$.} In the non-relativistic Skyrme phenomenology the latter is treated as an adjustable constant parameter with a typical value of $F_{so}(\rho) \equiv 3W_0/4 \approx 90\text{ MeV fm}^5$. It is worth noting that this value of the spin-orbit coupling strength is quite stable within the huge class of effective Skyrme forces considered in the literature \[11, 12\].

The basic framework for calculating the nuclear spin-orbit energy density functional $E_{so}[\rho, \vec{J}]$ is the density matrix expansion of Negele and Vautherin \[13\]. It generalizes the step-function distribution $\theta(k_f - |\vec{p}|)$ of nucleon momentum states for infinite nuclear matter to the situation of an inhomogeneous many-nucleon system with a local density $\rho(\vec{r})$ and a local spin-orbit density...
Fig. 1: Three-loop Hartree and Fock related to iterated one-pion exchange.

\[ \vec{J}(\vec{r}). \] Using this technique the Hartree diagrams of scalar and vector boson exchange between nucleons lead to the well-known expression for the spin-orbit coupling strength:

\[ F_{so}(\rho) = \frac{1}{4M^2(\rho)} \left( \frac{g_{\sigma N}^2}{m_\sigma^2} + \frac{g_{\omega N}^2}{m_\omega^2} \right) = \frac{G_S + G_V}{4M^2(\rho)}. \] (3)

Choosing typical values for the ratios of coupling constants to boson masses, \( G_S \simeq G_V \simeq 11 \text{ fm}^2 \), together with an effective (in-medium) nucleon mass of \( M^*(\rho_0) \simeq 0.73M_N \) one reproduces the empirical value of 90 MeV fm\(^5\) for the spin-orbit coupling strength.

Let us now turn to the "wrong-sign" spin-orbit interaction from the iterated one-pion exchange tensor force. The corresponding three-loop Hartree and Fock diagrams are shown in Fig. 1. Their contributions to the density dependent strength function \( F_{so}(k_f) \) have been evaluated analytically in Ref.\[14\]. For the purpose of illustration we reproduce here the expression for the dominant two-body Hartree contribution:

\[ F_{so}(\rho) = \frac{g_4^4 m_\pi M_N}{64\pi f_\pi^2} \left\{ \frac{1}{m_\pi^2 + 4k_f^2} - \frac{3}{8k_f^2} \ln \frac{m_\pi^2 + 4k_f^2}{m_\pi^2} \right\}, \] (4)

with the Fermi momentum \( k_f \) related to the nucleon density in the usual way: \( \rho = 2k_f^3/3\pi^2 \). When evaluated at nuclear matter saturation density, \( \rho_0 = 0.16 \text{ fm}^{-3} \), the expression in Eq.(4) leads to a large "wrong-sign" spin-orbit coupling strength of \( F_{so}(\rho_0) = -86.5 \text{ MeV fm}^5 \). Note that this term scales linearly with the large nucleon mass \( M_N = 939 \text{ MeV} \). This unfamiliar behavior of a spin-orbit coupling is obviously not a relativistic effect. It comes from the energy denominator of the iterated pion-exchange, which is proportional to the difference of small nucleon kinetic energies. As it stands, this large spin-orbit Hartree term (4) would simply cancel the scalar-vector mean field term of Eq.(3) at normal nuclear matter density.

The complete spin-orbit coupling strength from iterated 1\(\pi\)-exchange, with Fock term and Pauli blocking corrections included [14], is shown in Fig. 2 as a function of density. At nuclear matter saturation density \( \rho_0 \) a sizeable spin-orbit coupling strength \( F_{so}(\rho_0) = -47.4 \text{ MeV fm}^5 \), amounting to about one half of the empirical value but with the "wrong" sign, still remains.

We demonstrate now that this "wrong-sign" spin-orbit term is compensated by the three-body effects related to 2\(\pi\)-exchange with virtual \( \Delta(1232) \) isobar excitation. The dominant Hartree diagram in Fig. 3 leads to a spin-orbit strength function [15]:

\[ F_{so}(\rho) = \frac{g_4^4}{8\pi^2 \Delta f_\pi^4} \left\{ \frac{m_\pi^2 k_f + 2k_f^3}{m_\pi^2 + 4k_f^2} - \frac{m_\pi^2}{4k_f} \ln \frac{m_\pi^2 + 4k_f^2}{m_\pi^2} \right\}, \] (5)

with \( \Delta = 293 \text{ MeV} \) the delta-nucleon mass splitting which, notably, is a "small" scale just like the pion mass \( m_\pi \) and the nuclear Fermi momentum \( k_f \) when compared with the spontaneous chiral symmetry breaking scale, \( 4\pi f_\pi \simeq 1.2 \text{ GeV} \). Again, this contribution to the spin-orbit coupling is \emph{not} a relativistic effect. It originates primarily from the spin- and momentum dependence of the pion-baryon coupling and the fact that all three participating nucleons are from the
filled Fermi sea. Adding the very small contribution from the Fock diagrams [15], the spin-orbit coupling strength $F_{so}(\rho)$ shown in Fig. 4 results. As expected for a genuine three-body effect it grows with the nucleon density $\rho$. At nuclear matter saturation density one reads off the value $F_{so}(\rho_0) = 47.5$ MeV fm$^5$ which cancels the "wrong-sign" spin-orbit term from the second order $1\pi$-exchange tensor force. Further relativistic spin-orbit effects from the irreducible two-pion exchange (scaling as $1/M_N$ with the nucleon mass) have been investigated in Ref.[16]. These come out to be relatively small such that they do not affect in a significant way the cancellation mechanism between "wrong-sign" term from the iterated $1\pi$-exchange and "correct-sign" term from the Fujita-Miyazawa three-nucleon interaction. One might also consider $\rho_0/2$ (where the density-gradient is maximal) as the density relevant for estimating the spin-orbit coupling strength. At the same time, a detailed analysis of the spin-orbit strength function $F_{so}(\rho)$ in a chiral effective field theory approach [15] leads to a larger three-body contribution as shown by the dashed curve in Fig. 4. With this refined analysis the balance at $\rho_0/2 = 0.08$ fm$^{-3}$ goes as $-58.1$ MeV fm$^5$ against $+48.3$ MeV fm$^5$. This gives an impression of the uncertainties involved.

It has been pointed out that the strength of the short-range spin-orbit interaction, $F_{so}(\rho_0) \approx 90$ MeV fm$^5$ needed for (non-relativistic) nuclear structure calculations, is in perfect agreement with the one extracted from realistic nucleon-nucleon potentials (see Tables I and II in Ref.[16]). This intimate connection between the strength of the spin-orbit interaction in nuclei and free $NN$ scattering is further corroborated by the recent work of the Tübingen group [17]. Performing relativistic Brueckner calculations for the in-medium nucleon scalar and vector self-energies, they find that these strong mean-fields (of opposite sign) are almost entirely driven by the short-range spin-orbit part of the $NN$ potential used as input. If the corresponding strength parameter is tuned to zero both the Lorentz scalar and vector mean-fields tend to vanish completely (see Fig. 10 in Ref.[17]).

We thus conclude at this point that the nuclear spin-orbit interaction is subject to a subtle balance between three major pieces. The longer range spin-orbit term from the iterated (2nd or-
Fig. 3: Three-body Hartree and Fock diagrams of $2\pi$-exchange with virtual $\Delta(1232)$-isobar excitation.

Fig. 4: Spin-orbit coupling strength generated by the $2\pi$-exchange three-nucleon interaction. Full curve: Fujita-Miyazawa mechanism involving virtual $\Delta(1232)$-excitation. Dashed curve: Chiral effective field theory approach using the low-energy constant $c_3 = -3.9 \text{ GeV}^{-1}$ related to the nucleon axial polarizability [15].

der) pion exchange tensor force is canceled by the Fujita-Miyazawa type three-body contributions involving an intermediate $\Delta$ isobar. As a consequence the short-range scalar and vector pieces, familiar from Dirac-Hartree phenomenology, can account for the observed spin-orbit strength. Of course, an alternative book-keeping would also work equally well: one could have balanced the scalar-vector mean field contribution against just the Hartree term from iterated pion exchange and built up the spin-orbit strength by the Pauli blocking and Fock exchange pieces of in-medium second order pion exchange together with the three-body spin-orbit force. A detailed assessment of the interplay of all these different mechanisms becomes possible by comparing their actions in nuclei and $\Lambda$ hypernuclei.

3 Spin-orbit coupling for $\Lambda$ hyperons

The situation for $\Lambda$ hypernuclei is in a certain sense simpler since the aspect of self-consistency for the mean field potentials is not an issue in this case. A spherical mean field potential provided by the core nucleus determines already the single-particle motion of the $\Lambda$-hyperon. Systematic analyses by Millener, Dover and Gal [18] have shown that the empirical single-particle energies of
a Λ bound in hypernuclei are well described over a wide range in mass numbers by an attractive mean field potential of depth $U_\Lambda \simeq -28$ MeV, i.e. about half as strong as the one for nucleons in nuclei. On the other hand, the Λ-nucleus spin-orbit coupling is found to be extraordinarily weak. For example, recent precision measurements [19] of E1-transitions from $p$- to $s$-shell orbitals in $^{13}_Λ$C give a $p_{3/2} - p_{1/2}$ spin-orbit splitting of only $(152 \pm 65)$ keV, to be compared with a value of about 6 MeV in ordinary nuclei.

The empirical finding that the Λ-nucleus spin-orbit coupling appears to be negligibly small, in comparison with the strong spin-orbit interaction of nucleons in ordinary nuclei, posed an outstanding problem in low-energy hadronic physics. In relativistic scalar-vector mean field models a strong tensor coupling of the ω meson to the Λ hyperon, equal and of opposite sign to the vector coupling, was proposed as a possible solution [20], motivated by simple non-relativistic SU(6) quark model considerations in combination with the vector meson dominance hypothesis. We can demonstrate that there is actually a more natural source of cancellation in the hypernuclear many-body problem if one takes into account the "wrong-sign" spin-orbit coupling generated by the two-pion exchange with intermediate an Σ hyperon [21].

The pertinent quantity to extract the Λ-nuclear spin-orbit coupling is the spin-dependent part of the self-energy of a Λ hyperon interacting with weakly inhomogeneous nuclear matter. Let the Λ scatter from initial momentum $\vec{p} - \vec{q}/2$ to final momentum $\vec{p} + \vec{q}/2$. The spin-orbit part of the self-energy is then:

$$\Sigma_{\text{spin}} = \frac{i}{2} \vec{\sigma} \cdot (\vec{q} \times \vec{p}) U_{\Lambda ls}(\rho) ,$$

(6)

where the density-dependent spin-orbit strength $U_{\Lambda ls}(\rho)$ is taken in the limit of zero external momenta: $\vec{p} = \vec{q} = 0$. Its value at nuclear matter saturation density enters as a strength parameter in the more familiar spin-orbit Hamiltonian of the shell model:

$$\mathcal{H}_{\Lambda ls} = U_{\Lambda ls}(\rho_0) \frac{1}{2r} \frac{df(r)}{dr} \vec{\sigma} \cdot \vec{L} .$$

(7)

where $f(r)$ is a normalized nuclear density profile and $\vec{L} = \vec{r} \times \vec{p}$ the orbital angular momentum.

The key observation [21] is now that the iterated one-pion exchange with an intermediate Σ hyperon (see Fig. 5) also generates a sizeable Λ-nuclear spin-orbit coupling, with "wrong" sign opposite to that from scalar and vector mean fields. The basic mechanism behind it is again the spin- and momentum dependence of the pion-baryon interaction at second order. The prefactor $i \vec{\sigma} \times \vec{q}$ is immediately identified by rewriting the product of πΛΣ-interaction vertices $\vec{\sigma} \cdot (\vec{l} - \vec{q}/2) \vec{\sigma} \cdot (\vec{l} + \vec{q}/2)$, with $\vec{l}$ the momentum of the intermediate Σ hyperon. The other factor $\vec{p}$ emerges from the non-relativistic energy denominator which includes also the small ΣΛ-mass splitting $M_\Sigma - M_\Lambda = 77.5$ MeV.
Fig. 6: The spin-orbit coupling strength $U_{\Lambda ls}(\rho)^{(2\pi \Sigma)}$ of a $\Lambda$ hyperon as function of the nuclear density $\rho$. The lower curve shows the long-range contribution from iterated pion exchange with a $\Sigma$ hyperon in the intermediate state. The two upper curves include in addition the short-range contribution, $U_{\Lambda ls}(\rho)^{(sr)} = 24.8 C_{ls} \text{MeV fm}^2 \cdot \rho/\rho_0$, with $C_{ls} = 2/3$ and $1/2$ as options.

The lower curve in Fig. 6 shows the $\Lambda$-nuclear spin-orbit coupling strength $U_{\Lambda ls}(\rho)^{(2\pi \Sigma)}$ resulting from these long-range two-pion exchange processes as a function of the nucleon density (for analytical expressions, see Eqs.(11-15) in Ref.[21]). With a value of $U_{\Lambda ls}(\rho_0) \simeq -15 \text{MeV fm}^2$ at normal nuclear matter density $\rho_0 = 0.16 \text{fm}^{-3}$ it is comparable in magnitude but of opposite sign with respect to the spin-orbit terms from scalar and vector mean fields. The upper two curves in Fig. 6 include in addition a short-range component of the $\Lambda$-nuclear spin-orbit coupling:

$$U_{\Lambda ls}(\rho)^{(sr)} = C_{ls} \frac{M_N^2}{M_\Lambda^2} U_{Nls}(\rho)^{(sr)}.$$

which has been scaled to the one of nucleons. This coefficient $C_{ls}$ parameterizes the ratio of relevant coupling strengths. We have varied it between $C_{ls} = 2/3$ (following from naive quark model considerations) and $C_{ls} = 1/2$ as suggested by in-medium QCD sum rule calculations of the Lorentz scalar and vector mean fields [22]. For the nucleonic spin-orbit coupling strength we take the value $U_{Nls}(k_{f_0})^{(sr)} = 3\rho_0 W_0/2 \simeq 30 \text{MeV fm}^2$ from shell model calculations. One observes from Fig. 6 an almost complete cancellation between short-range and long-range contributions. This balance offers a natural explanation of the empirically observed small spin-orbit splitting in $\Lambda$ hypernuclei.

It is important to note that the nuclear three-body contributions which compensated the "wrong-sign" spin-orbit terms (from the second order pion exchange tensor force) does not exist in hypernuclei. The absence of an analogous three-body mechanism for a $\Lambda$ in the hypernucleus becomes immediately clear by inspection of Fig. 7. Replacing the external nucleon by a $\Lambda$ introduces an intermediate $\Sigma$ hyperon. However, since there is no filled Fermi sea for hyperons, a three-body force analogous to that of Fig. 7 does not exist in this case. Moreover, it has recently been shown in Ref.[23] that the cancellation mechanism proposed in Ref.[21] is not disturbed by the inclusion of analogous $2\pi$-exchange processes with all relevant decuplet baryons ($\Delta(1232)$ and $\Sigma^*(1385)$) in the intermediate state. These effects have alternating signs from spin-sums and are suppressed by considerably larger mass-splittings in the energy denominator.
The emerging picture of the nuclear and hypernuclear spin-orbit interaction is an intriguing one. The spin-orbit interaction of nucleons is predominantly of short range because the longer range $2\pi$-exchange components find a mechanism of self-cancellation involving an important three-body term. The smallness of the $\Lambda$-nuclear spin-orbit coupling, on the other hand, reveals the existence of long range $2\pi$-exchange component of the "wrong sign" which balances the short-distance mechanisms. In a recent paper [24] it has been demonstrated that this scenario works indeed for actual (finite) $\Lambda$ hypernuclei.

4 Summary and conclusions

Let us summarize our findings. It is well known that large Lorentz scalar and vector mean fields of opposite sign can generate (via the lower components of Dirac spinors) the strong spin-orbit coupling of nucleons in nuclei as a relativistic effect. A successful nuclear structure phenomenology has been developed on the basis of this mechanism. Less well known is the fact that two-pion exchange also produces sizeable spin-orbit coupling terms. The mechanism behind these is quite different and not of relativistic origin. These spin-orbit couplings arise directly from the spin- and momentum dependence of the pion-nucleon interaction together with small energy denominators which enhance the effects.

When working out these long range spin-orbit coupling terms (which depend only on well-known hadronic parameters) one finds that a "wrong-sign" contribution from the second order pion-exchange tensor force gets canceled by a contribution from the three-nucleon interaction of Fujita-Miyazawa type (mediated by $2\pi$-exchange with virtual $\Delta$ isobar excitation). We have explicitly verified this cancellation mechanism by studying the spin-orbit strength function $F_{so}(\rho)$ in the nuclear energy density functional around nuclear matter saturation density, $\rho_0 = 0.16 \text{fm}^{-3}$.

When searching for the microscopic origin of the strong Lorentz scalar and vector mean fields one finds, within relativistic Brueckner calculations, that they are mainly caused by the short-range spin-orbit part of the $NN$ potential used as an input. The same connection is observed when comparing directly the spin-orbit interaction strength $3W_0/4$ in the Skyrme phenomenology with the one extracted from realistic nucleon-nucleon potentials. An alternative approach to the strong Lorentz scalar and vector mean fields in nuclear matter is provided by QCD sum rules which link those fields to changes of quark condensates at finite baryon density.

For a $\Lambda$ in a hypernucleus the balance between short-range and long-range components of the spin-orbit coupling is qualitatively different as compared to ordinary nuclei. The three-
body contribution (induced by $2\pi$-exchange) is now absent. The (scalar-vector) mean field term is canceled by a "wrong-sign" contribution from iterated pion-exchange with an intermediate $\Sigma$ hyperon. The small mass splitting $M_\Sigma - M_\Lambda = 77.5\,\text{MeV}$ is crucial in order to make this cancellation so effective. We note in passing that a similar cancellation mechanism is expected to be at work for the $\Sigma$-nuclear spin-orbit interaction [25]. Unfortunately, the prospects for observing this effect are poor because of the recently established repulsive nature of the $\Sigma$-nuclear bulk potential [26].

This description of spin-orbit interactions guided by chiral effective field theory, with a short-range (mean field) component and long-range contributions of alternating signs from $2\pi$-exchange, can explain (at least qualitatively) the pronounced difference in strength between the spin-orbit interaction of nucleons in ordinary nuclei and a $\Lambda$ in hypernuclei.

Acknowledgements
We thank Avraham Gal for many stimulating discussions.

References
[1] B.D. Serot and J.D. Walecka, *Int. J. Mod. Phys.* E6 (1997) 515; and references therein.
[2] P. Ring, *Prog. Part. Nucl. Phys.* 37 (1996) 193; and references therein.
[3] H.P. Dürr, *Phys. Rev.* 103 (1956) 469.
[4] L.D. Miller, *Ann. of Phys.* 91 (1975) 40.
[5] R. Brockmann and W. Weise, *Phys. Rev.* C 16 (1977) 1282.
[6] P. Finelli, N. Kaiser, D. Vretenar and W. Weise, *Eur. Phys. J.* A17 (2003) 573; *Nucl. Phys. A* 735 (2004) 449; *Nucl. Phys. A* 770 (2006) 1.
[7] N. Kaiser, R. Brockmann and W. Weise, *Nucl. Phys. A* 625 (1997) 758.
[8] S.C. Pieper and V.R. Pandharipande, *Phys. Rev. Lett.* 70 (1993) 2541.
[9] S.C. Pieper, V.R. Pandharipande, R.B. Wiringa, and J. Carlson, *Phys. Rev.* C64 (2001) 014001; and references therein.
[10] J. Fujita and H. Miyazawa, *Prog. Theor. Phys.* 17 (1957) 360; 17 (1957) 366.
[11] M. Bender, P.-H. Heenen and P.-G. Reinhard, *Rev. Mod. Phys.* 75 (2003) 121; and references therein.
[12] E. Chabanat, P. Bonche, P. Haensel, J. Meyer, and R. Schaeffer, *Nucl. Phys.* A635 (1998) 231.
[13] J.W. Negele and D. Vautherin, *Phys. Rev.* C5 (1972) 1472.
[14] N. Kaiser, S. Fritsch and W. Weise, *Nucl. Phys.* A724 (2003) 47; and references therein.
[15] N. Kaiser, *Phys. Rev.* C68 (2003) 054001.
[16] N. Kaiser, *Phys. Rev.* C70 (2004) 034307.
[17] O. Plohl and C. Fuchs, *Phys. Rev.* C74 (2006) 034325.
[18] D.J. Millener, C.B. Dover, and A. Gal, *Phys. Rev.* C38 (1988) 2700.
[19] S. Ajimura et al., *Phys. Rev. Lett.* 86 (2001) 4255.
[20] B.K. Jennings, *Phys. Lett.* B246 (1990) 325.
[21] N. Kaiser and W. Weise, *Phys. Rev.* C71 (2005) 015203.
[22] X. Jin and R.J. Furnstahl, *Phys. Rev.* C49 (1994) 1190.
[23] J. Martin Camalich and M.J. Vicente Vacas, *Phys. Rev.* C76 (2007) 068201.
[24] P. Finelli, N. Kaiser, D. Vretenar and W. Weise, *Phys. Lett.* B658 (2007) 90.
[25] N. Kaiser, *Phys. Rev.* C76 (2007) 068201.
[26] E. Friedman and A. Gal, *Phys. Reports* 452 (2007) 89.