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Dirac Phenomenology and Hyperon-Nucleus Interactions

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We discuss various aspects of hyperon-nucleus interactions in the relativistic mean field theory. First, characteristics of Λ , Σ , and Ξ hypernuclei, as well as multi-strange baryonic objects, are investigated. The spin-orbit splittings and magnetic moments are shown to be very sensitive to the value of the tensor coupling $f_{\omega\gamma}$. Second, optical potentials for Λ and Σ scattering off nuclei are developed based on a global nucleon-nucleus Dirac optical potential and $SU(3)$ symmetry. The tensor coupling has a large effect on the predictions for the analyzing power. Third, the Dirac approach is used in the calculations of the non-mesonic decay of Λ hypernuclei. The large discrepancy between the decay rates and data suggests the need for additional meson exchanges.

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A hyperon embedded in the nuclear medium is an unique probe of the deep nuclear interior[1]. Hypernuclei, as bound systems of different types of baryons, represent a sound generalization of traditional nuclear matter to a many-body baryonic system. Their study provides a direct test of various models for the baryon-baryon interaction[2-5]. Moreover, due to the strangeness contents of their constituent(s), hypernuclei may be used for investigating the possible manifestation of the sub-hadronic degrees of freedom in nuclear medium (for instance via studying the Pauli blocking effects).

Many studies in the last decade have proved the applicability of the relativistic mean field theory (RMFT)[6, 7] to the description of nuclear structure as well as nuclear dynamics. The success of RMFT in the nuclear case inspired quite naturally attempts to describe, within the same theory, the more general baryon-nucleus interaction[8-16]. The consistent description of both nuclear and hypernuclear systems[17-20] can be regarded as a great success, putting confidence in the relativistic mean field approach.

Much effort has been devoted to searching for a direct test of genuine relativistic dynamics. The study of hypernuclear currents and magnetic moments in particular has been regarded to be the region where the relativistic and nonrelativistic models differ in their predictions[21, 22]. However, the most recent calculations[23, 24] revealed that it is hard to expect an unambiguous experimental signature.

Multi-strange objects, either baryonic systems or strange quark matter have been the subject of many theoretical studies[2, 11-15, 25]. They are of much interest for the physics of relativistic heavy ion collisions[26] and for astrophysical considerations[27] although the experimental information is limited to the ground state binding energies of three double- Λ hypernuclei: $^8_{\Lambda\Lambda}He$, $^{10}_{\Lambda\Lambda}Be$ and $^{13}_{\Lambda\Lambda}B$ [28].

Hyperon-nucleus scattering[29, 30] is still a rather unexplored field of hyperon physics. Currently there is no experimental data available, however some information on Σ^+ scattering off ^{12}C can soon be expected from KEK experiments[31]. Present calculations thus serve as a first estimate of the hyperon-nucleus scattering observables.

Dirac phenomenology can also be used in the calculation of non-mesonic decay of Λ hypernuclei[32, 33]. This uses both the bound state wave function of the initial state and the scattering wave function of the two out going nucleons.

In the next section, we introduce the underlying RMFT model (subsection A) and discuss the choice of parameters (subsection B). Application to the spectroscopy of Λ , Σ and Ξ hypernuclei is presented in Section III(A), while Section III(B) is devoted to the hypernuclear currents and magnetic moments. A brief summary of the RMFT predictions for the multi-strange baryonic systems is given in Section IV. In Section V, we apply Dirac phenomenology to the description of a hyperon-nucleus scattering. The non-mesonic decay of Λ hypernuclei is discussed in Section VI. Conclusions are drawn in Section VI.

II. RELATIVISTIC MEAN FIELD MODEL

A. Lagrangian

The framework for our investigations is the relativistic mean field theory (RMFT)[6]. It describes a nucleus as a system of Dirac spinors (nucleons) interacting via meson fields in the mean field approximation. Our aim is to calculate hypernuclear systems and so we extend the original model for ordinary nuclei by the strange particle (Λ , Σ , Ξ) sector. We start from the Lagrangian density of the form

$$\mathcal{L} = \mathcal{L}_N + \mathcal{L}_Y, \quad Y = \Lambda, \Sigma, \Xi, \quad (1)$$

$$\mathcal{L}_Y = \bar{\Psi}_Y [i \gamma_\mu \partial^\mu - g_{\omega Y} \gamma_\mu V^\mu - (M_Y + g_{\sigma Y} \phi)] \Psi_Y + \mathcal{L}_N + \mathcal{L}_\pi + \mathcal{L}_T. \quad (2)$$

In Eq. (1), $\mathcal{L}_N$ denotes the standard Lagrangian of the mean field theory for nuclei. Its detailed form can be found elsewhere[6, 7]. $\mathcal{L}_N$ involves nucleons (Ψ_N), scalar σ mesons (ϕ), vector ω mesons (V^μ), vector isovector ρ mesons ($\vec{\rho}^\mu$) and the photon (A^μ) (for the Coulomb field). In the case of the nonlinear model the scalar meson self-couplings are included as well:

$$\mathcal{L}_{\pi 1} = -\frac{1}{3} g_2 \phi^3 - \frac{1}{4} g_3 \phi^4. \quad (3)$$

In general, the $\mathcal{L}_Y$ contains contributions $\mathcal{L}_{\rho Y}$ and $\mathcal{L}_{A Y}$ from the interactions of a hyperon with the ρ meson and Coulomb field, respectively. For a particular hyperon they acquire the following form:

$$\mathcal{L}_{\rho \Lambda} = \mathcal{L}_{A \Lambda} = 0, \quad (4)$$

$$\mathcal{L}_{\rho \Xi} + \mathcal{L}_{A \Xi} = -\bar{\Psi}_\Xi \left[\frac{1}{2} g_{\rho \Xi} \vec{\tau}_\Xi \cdot \vec{\rho}^\mu \gamma_\mu + \frac{e}{2} (\tau_{3\Xi} - 1) \gamma_\mu A^\mu \right] \Psi_\Xi, \quad (5)$$

$$\mathcal{L}_{\rho \Sigma} + \mathcal{L}_{A \Sigma} = -\bar{\Psi}_\Sigma \left(\frac{g_{\rho \Sigma}}{2} \gamma_\mu \Theta_{\Sigma k}^\mu + \frac{e}{2} \gamma_\mu A^\mu (\tau_{3\Sigma})_{kk} \right) \Psi_{\Sigma k}, \quad (6)$$

where

$$\Sigma_{ij} = \begin{pmatrix} \Psi_{\Sigma^0} & \sqrt{2} \Psi_{\Sigma^+} \\ \sqrt{2} \Psi_{\Sigma^-} & -\Psi_{\Sigma^0} \end{pmatrix}, \quad (7)$$

and

$$\Theta_{ij}^\mu = \begin{pmatrix} \rho_0^\mu & \sqrt{2} \rho_+^\mu \\ \sqrt{2} \rho_-^\mu & -\rho_0^\mu \end{pmatrix}. \quad (8)$$

Finally, we included the Lagrangian density $\mathcal{L}_T$ that describes the ω - Y anomalous coupling term,

$$\mathcal{L}_T = \frac{f_{\omega Y}}{2M_Y} \bar{\Psi}_Y \sigma^{\mu\nu} \partial_\nu V_\mu \Psi_Y. \quad (9)$$

In the following we neglect the contributions of anti(quasi)particles (so called "no-sea" approximation) and quantum fluctuations of meson fields ("mean-field" approximation). Using standard variational techniques we arrive at a system of the Dirac and Klein-Gordon equations for baryons and meson fields, respectively. In the static limit they acquire the form:

$$[-\vec{\alpha}(\vec{\nabla} + g_{\omega i} \vec{V}) + g_{\sigma i} V_0 + \beta(M_i + g_{\sigma i} \phi) - \delta_{Yi} i \frac{f_{\omega Y}}{2M_Y} \beta \vec{\alpha} \cdot \vec{\nabla} V_0 +$$

$$I_{3i} g_{\rho i} \vec{b}_0 + Q_i e A_0] \psi_{\alpha i} = \epsilon_{\alpha i} \psi_{\alpha i}, \quad i = N, \Lambda, \Sigma, \Xi, \quad (10)$$

$$(-\Delta + m_j^2) B_j = \sum_i g_{ji} \rho_{ji} - \delta_{Yj} \delta_{j\omega} \frac{f_{\omega Y}}{2M_Y} \rho_T + \delta_{j\sigma} (-g_2 \phi^2 - g_3 \phi^3), \quad (11)$$

$$-\Delta A_0 = e \sum_i \rho_{\alpha i}, \quad (12)$$

where

$$B_\sigma = \phi, \quad B_{\omega 0} = V_0, \quad B_{\vec{\sigma}} = \vec{V}, \quad B_{\rho_0} = b_0,$$

and I_{3i} and Q_i are the third isospin component and the charge of the baryon, respectively.

The densities in Eqs.(11) and (12) are defined as:

$$\rho_{ji} = \sum_{\alpha} \bar{\psi}_{\alpha i} D_j \psi_{\alpha i}, \quad j = \sigma, \omega_0, \vec{\omega}, \rho_0,$$

$$\rho_{\alpha i} = \sum_{\alpha} \bar{\psi}_{\alpha i} Q_i \psi_{\alpha i},$$

$$\rho_T = \vec{\nabla} \cdot \sum_{\alpha} \bar{\psi}_{\alpha i} i \vec{\alpha} \psi_{\alpha i},$$

$$D_\sigma = -\vec{\tau}_i, \quad D_{\omega_0} = \gamma_0, \quad D_{\vec{\sigma}} = \vec{\tau}_i, \quad D_{\rho_0} = \gamma_0 I_{3i},$$

where in summations α runs over all occupied particle states.

Hypernuclei studied here are systems with one hyperon added to a spherical (closed shell) nuclear core. Such systems with an odd number of particles are no longer time-reversal invariant[34, 35]. As a consequence, the spatial components of vector fields and baryon currents in Eqs. (10)-(11) do not vanish. They have been

omitted in the majority of the hypernuclear RMFT studies and, indeed, they are of little importance in spectroscopic calculations of nuclear systems[36, 37]. On the other hand, retaining the spatial components $\vec{V}$ in the Hartree calculations of the entire “closed core+1” system is critical for a consistent inclusion of the effect of the nuclear core response[24] and consequently for a description of the currents and the magnetic moments[21, 22, 37], in particular.

Contributions of the three-vector $\vec{\rho}$ and $\vec{A}$ fields are negligible[22] and so they are not included. The tensor coupling term for nucleons is omitted as well since the $f_{\omega N}$ coupling constant is small. Finally, we neglect the ρYY tensor coupling terms whose effect was also found to be very small[20].

The system of coupled field Eqs. (10)-(12) for both baryons (N, Y) and meson fields ϕ , V_0 , $\vec{V}$, b_0 , A_0 is solved fully self-consistently. This is important particularly for calculating the hypernuclear magnetic moments and for evaluation of the isovector ρ contribution to the hyperon binding energies as will be seen in the following sections.

B. Parametrization

The parameters of the nuclear RMFT models are usually fitted to the properties of spherical nuclei[7, 38, 39]. In contrast to the nuclear matter fits[6] this approach determines also scalar meson mass m_σ and coupling constant $g_{\sigma N}$. The calculations presented in the following sections were done using the nonlinear parameter set *SH* of Sharma *et al.*[40] and linear parametrization ($g_2 = g_3 = 0$) *L* of Horowitz and Serot[38] for the nucleonic sector. We chose the nonlinear model *SH* since it is numerically more stable than the widely used *NL*[39] which already shows instabilities for more dense nuclear systems, as e.g. ^{12}C [7]. The masses and coupling constants of the *L* and *SH* parametrizations are listed in Table I.

The ratios $\alpha_{\sigma A} = g_{\sigma A}/g_{\sigma N}$, $\alpha_{\omega A} = g_{\omega A}/g_{\omega N}$ for the Λ coupling are usually adjusted[11, 12, 14] to the recent (π^+ , K^+) data[41]. The Λ couplings that describe the Λ single particle states reasonably well are located along a line in the $(\alpha_{\sigma A}, \alpha_{\omega A})$ plane[12-14]. This linear dependence results from the constraint to obtain the correct well depth U_Λ of ≈ 30 MeV in nuclear matter[15]:

$$U_\Lambda = -(\alpha_{\sigma A} S_N + \alpha_{\omega A} V_N) = M_N \left(1 - \frac{M_N^2}{M_N^2}\right) \alpha_{\sigma A} - \frac{g_{\omega N}^2}{m_\omega^2 \rho_0} \alpha_{\omega A},$$

where S_N and V_N are the scalar and vector potentials for nucleon, respectively. In many related works[8, 11, 14] the tensor coupling contribution $\mathcal{L}_T$ in Eq. (2) has not been considered. The experimentally measured small spin orbit splitting has been achieved by choosing weak Λ coupling ($\alpha_{\sigma A} \approx \frac{1}{3}$). The resulting spin orbit potential which there goes like: $\frac{1}{M_N^2} (\vec{V}(\alpha_{\sigma A} S_N - \alpha_{\omega A} V_N))$ is then reduced by $\approx \frac{1}{3} (M_N^2)^2 \approx 0.1$ relatively to nucleon case.

Unfortunately, it is hard to explain such a large SU(3) symmetry breaking[17]. Moreover, it is impossible to reconcile such weak hyperon coupling with the quark

model predictions[42] of a sizable spin orbit interaction for Σ hyperon. Finally, stronger Λ coupling, namely $\alpha_{\sigma A} \gtrsim 0.5$, is required in order to get bounds on a neutron star mass[27] compatible with the observations.

Recently it was found[17, 18] that a consistent description of hypernuclear data can be achieved even with larger values of the meson couplings (in agreement with SU(3)) if the tensor coupling term $\mathcal{L}_T$ (9) is taken into account in the Lagrangian $\mathcal{L}_Y$ (3).

For the coupling of the ω meson to the hyperons we have used the naive quark model values of the ratio to the ω -nucleon coupling. We have adopted the values of ωYY couplings directly from Dover and Gal[3]. The ρ couplings can be easily determined by the same techniques using SU(3) Clebsch-Gordan coefficients.

For illustration of the effect of the ωYY tensor coupling to the hypernuclear spin-orbit splitting, we adopted for $f_{\omega Y}$ also values from Nijmegen models F[43] and D[44] and values suggested by Cohen and Weber[18].

While we used the quark model to determine the values of coupling constants $g_{\omega Y}$, $f_{\omega Y}$ and $g_{\rho Y}$, the $g_{\sigma Y}$ couplings were deduced from the available experimental information on hyperon binding in the nuclear medium. The hypernuclear levels are very sensitive to the precise value of the couplings due to large cancellations between scalar and vector potentials and therefore the coupling adjustment is required for a quantitative agreement with the hypernuclear levels. The SU(3) inspired values of $g_{\sigma Y}$ and $g_{\rho Y}$ are denoted by P3 in Table II.

In order to study the dependence on Λ couplings in some particular cases we used as well other pairs (P1-P3) of ratios $\alpha_{\sigma A}$, $\alpha_{\omega A}$.

Due to insufficient experimental information on Σ and Ξ hypernuclei, we chose $g_{\sigma Y}$ to roughly obtain the estimated Σ (Ξ) potential depth in nuclear matter. Analyses of the level shifts in Σ^- atoms[45, 46] lead to a potential depth $U_\Sigma \approx 20 - 30$ MeV for a Σ imbedded in nuclear matter. A similar value $U_\Xi \approx 20 - 25$ MeV is compatible with a scarce data from Ξ^- hypernuclei observed in emulsion experiments with K^- beams[47]. The values of $g_{\rho Y}$ presented in Table II give $U_Y \approx 25 - 30$ MeV.

The values of $\alpha_{\sigma Y} = g_{\sigma Y}/g_{\sigma N}$, $\alpha_{\omega Y} = g_{\omega Y}/g_{\omega N}$, $\alpha_{\rho Y} = g_{\rho Y}/g_{\rho N}$ and $\alpha_{\tau Y} = f_{\omega Y}/g_{\omega Y}$ for all the models used in the following sections are presented in Table II. We have neglected the finite width of the Σ hyperon states due to $\Sigma N \rightarrow \Lambda N$ conversion and $\Sigma N \rightarrow \Sigma N$ charge exchange. The question of the width of the Σ states is still open and goes beyond the scope of this work.

III. Λ , Σ AND Ξ HYPERNUCLEI

A. Shell model for hyperons

The RMFT models introduced in Section 2 give reasonable description of hypernuclear characteristics – hyperon binding in nuclear matter, spin orbit interaction, single particle energies. In this subsection we illustrate this statement on several examples.

In Fig. 1 we present the Λ single particle energies in hypernuclei with $A = 12 - 209$ calculated using model *SH* from Table I, parameter set P3 from Table II

and $f_{\Lambda\Lambda}/g_{\omega\Lambda} = -1$. The figure shows a good agreement (within experimental errors ≈ 1 MeV) of the RMFT results with the data from (π^+, K^+) production[41]. The evolution of binding energies of a neutral Σ^0 is similar to that of Λ in Fig. 1 as long as we assume the same strength of Λ and Σ coupling in nuclear medium.

Figure 2 demonstrates the importance of the Coulomb and isovector (ρ meson) interactions in $\Sigma^\pm$ hypernuclei. The following discussion is valid however for Ξ hypernuclei, as well[20]. The attractive Coulomb interaction (Fig. 2(a)) leads to a considerably stronger binding of Σ^- (Ξ^-) in nuclear medium when compared with neutral hyperons Λ , Σ^0 and Ξ^0 . This allows the population of highly excited hypernuclear states[16]. For Σ^+ the repulsive Coulomb potential decreases the binding of the Σ^+ hyperon (Fig. 2(b)). The effect grows with charge number Z leading even to decreasing binding B_{Σ^+} when going from Ca to Pb . Consequently, only s and p bound states were found in Σ^+ hypernuclei for $\alpha_{\Sigma} = 0$.

However, the $\rho Y Y$ coupling affects the previous results appreciably as can be seen also in Fig. 2 (solid line). The self consistent calculations[20] of the entire "core + hyperon" systems predict slightly attractive (and thus nonzero!) ρ meson contribution even for hypernuclei with $N = Z$ core (O, Ca). In Zr and Pb the ρ contribution is repulsive (attractive) for both Σ^-, Ξ^- (Σ^+, Ξ^0) and thus compensates to some extent the effect of the attractive (repulsive) Coulomb potential. In particular, the binding B_{Σ^+} converts from decreasing to an increasing function of A in the systems with a neutron excess. This result strongly depends on the value of the $g_{\rho Y}$ coupling constant; in Ref. [16] where $\alpha_{\Sigma} = 0.6$ was used, no such conversion of the slope was observed.

The tensor coupling (9) is crucial for evaluating the hypernuclear spin orbit interaction. In Fig. 3, $VO(Y = \Lambda, \Sigma^0, \Xi^0)$ calculated using various tensor couplings α_{TY} from Table II serves as an example of quite different evolution of the spin orbit splitting for three kinds of hyperons. The Λ (Σ) hypernuclear spin orbit splitting is gradually decreasing (increasing) due to a negative (positive) α_{TY} . Extremely large negative values of $\alpha_{T\Sigma}$ result even in a change of the level ordering in Ξ hypernuclei. The above results become apparent from Schroedinger equivalent spin orbit potential[48]:

$$V_{\alpha}^Y \vec{l} \cdot \vec{s} = \frac{1}{2M_{eff}^2} \left[\frac{1}{r}, \left(g_{\omega Y} V_0' - g_{\sigma Y} \phi' + 2f_{\omega Y} \frac{M_{eff} l V_0'}{M_Y} \right) \right] \vec{l} \cdot \vec{s},$$

$$M_{eff} = M_Y - \frac{1}{2} (g_{\omega Y} V_0 - g_{\sigma Y} \phi).$$

For $f_{\omega Y} = 0$ the spin-orbit splitting (V_{α}^Y) for hyperons is reduced when compared to that one for nucleons due to a larger mass M_{eff} in the denominator and due to the smaller couplings to σ and ω mesons. The quark model values of the tensor coupling $f_{\omega Y}$ for Λ , Σ and Ξ hyperons differ in their strengths and signs and hence their contribution to the spin-orbit term is different. This contribution is comparable in magnitude with the original $\sigma - \omega$ part and is negative (positive) for Λ (Σ). It is negative and even larger than $\sigma - \omega$ term for Ξ . As a consequence, the spin orbit interaction for Λ nearly vanishes, it is almost doubled for Σ , and changes sign for Ξ .

It should be noted, however, that although the effects of tensor coupling are relatively large, the absolute shifts in energy levels amount to less than 1 MeV and this is still beyond the reach of contemporary experimental resolution[41].

B. Λ hypernuclear currents and magnetic moments

The relativistic and nonrelativistic models yield similar predictions for the isoscalar nuclear magnetic moments[49]. While the nonrelativistic (Schmidt) values are obtained directly in a single-particle approach, in the RMFT calculations the result comes from a compensation of two effects – the enhancement of the valence particle current due to the reduction of the effective nucleon mass and the contribution of the additional current from the polarized core[50]. This fundamental difference was expected to have a crucial consequence for Λ hypernuclei where the enhanced valence hyperon current was not entirely cancelled by core corrections due to different masses and couplings of Λ and N .

The first self consistent calculations[21, 22] of hypernuclear currents and magnetic moments did not include the tensor coupling contribution (3). While Cohen and Furnstahl[21] evaluated the effect of the linear core response to the hyperon in the local density approximation, Mareš and Žofka[22] made self consistent calculations of the entire "core + Λ " system and incorporated this effect automatically. Both calculations revealed remarkable differences between relativistic and Schmidt values.

The magnetic moments of several hypernuclei calculated in RMFT model L from Table I and parametrizations P1–P3 from Table II, with $f_{\omega\Lambda} = 0$ are presented in Table III. Since the Λ is a neutral particle the Schmidt value comes entirely from its anomalous moment. The deviations are then caused by nonvanishing contribution from polarized core. It is to be noted that the spatial components of the vector field $\vec{V}$ are crucial here – when they are omitted the magnetic moments agree with the nonrelativistic predictions. The differences between RMFT and Schmidt values are proportional to the $\alpha_{\omega\Lambda}$ and increase with A [22]. It should be stressed again that the above RMFT results were obtained when the anomalous $\omega\Lambda\Lambda$ coupling was not included.

Introducing the tensor coupling term (3) into the Lagrangian (2) influences the predictions of the baryon currents and hypernuclear magnetic moments considerably[23, 24]. An excellent discussion with analytical derivations of the previous numerical results[23] was given by Cohen and Noble[24], so we limit ourselves to a brief statement of the conclusions. The vector current $\bar{\Psi}_\Lambda^\dagger \vec{\alpha} \Psi_\Lambda$ is modified due to the tensor coupling term by the additional contribution: $\frac{\alpha_{T\Lambda}}{2M_\Lambda} \vec{V} \times (\bar{\Psi}_\Lambda^\dagger \vec{\rho} \vec{\Sigma} \Psi_\Lambda)$. The hypernuclear magnetic moment

$$\vec{\mu}_{HY} = \frac{1}{2} \int d^3\vec{r} \vec{r} \times \{ \vec{j}_\Lambda(\vec{r}) + \vec{j}_{core}(\vec{r}) \},$$

has contributions from the hyperon and core electromagnetic currents ($\vec{j}_\Lambda(\vec{r})$ comes from the anomalous moment μ_Λ). The Λ contribution to the hypernuclear magnetic

moment is close to the Schmidt value again. Evaluations of the core contribution $\mu_{core}[24]$ (which was the source of the previous deviations[21, 22] from the Schmidt value) lead for $f_{\Lambda}/g_{\Lambda} = -1$ and Λ hyperon in s state to $\mu_{core} \approx 0$. Consequently the μ_{HY} is predicted to be close to the Schmidt value. The nonvanishing core contribution predicted in Refs. [21, 22] remains for the Λ in states with higher orbital angular momentum. However, the measurements of the magnetic moments for these states are not anticipated in the near future.

The RMFT model with the $\omega Y Y$ tensor coupling leads to different conclusions concerning the detectable relativistic effects in μ_{HY} than previous calculations that have neglected this term. Reliable calculations of hypernuclear magnetic moments with careful evaluation of various effects have become important since their measurement is recently under consideration[51].

IV. MULTI-STRANGE BARYONIC SYSTEMS

The baryonic systems, with larger contents of strangeness, have already been subject to numerous studies[2, 11–15, 52] despite the fact that the only experimental evidence of the coexistence of hyperons in a nuclear matter is a few events of double- Λ hypernuclei ${}^6_{\Lambda\Lambda}He$, ${}^{10}_{\Lambda\Lambda}Be$, ${}^{13}_{\Lambda\Lambda}B$ [28]. Baryonic systems with strangeness $S \leq -3$ or their more strongly bound quark counterparts, “strangelets”[25], are expected to be produced in the relativistic heavy ion collisions[26]. Enhanced interaction radii of these objects lead to behaviour analogous to the “anomalons”[11]. It has long been recognized that a large number of strange particles will be present in the cores of neutron stars[25, 27].

The investigations of the bound systems with 2 or more hyperons is expected to provide additional information about the baryon-baryon interaction. This can help to distinguish between different interaction models and theoretical approaches, that are more or less equivalent in describing the $Y N$ scattering or single hypernuclei.

When applied to double- Λ hypernuclei the RMFT model, as formulated in Section 2, fails to incorporate correctly the contribution from the $\Lambda\Lambda$ interaction[53, 54]. The pertinent observables that contain the information about the $\Lambda\Lambda$ interaction are:

$$B_{\Lambda\Lambda}({}^A_{\Lambda\Lambda}Z) = B_{\Lambda}({}^A_{\Lambda}Z) + B_{\Lambda}({}^{A-1}_{\Lambda}Z),$$

$$\Delta B_{\Lambda\Lambda}({}^A_{\Lambda\Lambda}Z) = B_{\Lambda}({}^A_{\Lambda\Lambda}Z) - 2B_{\Lambda}({}^{A-1}_{\Lambda}Z).$$

The analysis of the data yields $\Delta B_{\Lambda\Lambda} \approx 4-5$ MeV, while in the RMFT calculations $\Delta B_{\Lambda\Lambda}$ is much smaller ($\Delta B_{\Lambda\Lambda} \leq 1$ MeV) and even decrease for stronger Λ couplings[15, 54]. It is worth mentioning that in the RMFT model presented above the $\Lambda\Lambda$ contribution comes from the relatively small changes of the meson fields due to adding the second hyperon. Recently, the problem of a small $\Lambda\Lambda$ binding was solved by Schaffner *et al.*, [15] by introducing an additional $\Lambda\Lambda$ interaction mediated by two mesons (scalar f_0 and vector ϕ) which couple exclusively to the strange baryons.

The RMFT calculations revealed a possibility of forming bound systems with an appreciable number of hyperons. The dependence of the binding energies per

particle, density distributions and rms radii of such systems on a number of hyperons n_{Λ} , results from a delicate interplay between the effect of Pauli blocking (Λ distinguishable from N) and the weaker ΛN interaction compared to the NN one. Different strengths of the Λ couplings lead to a qualitatively similar n_{Λ} dependence of the studied characteristics. Nevertheless, the use of weaker Λ couplings results in a shift of the hyperon density distribution towards the centre of the system and consequently to the smaller Λ rms radius[12]. The effect of the anomalous $\omega\Lambda\Lambda$ coupling (3) on the investigated characteristics is less than 3%[54], and therefore, it is not necessary to include the tensor coupling terms when one deals with the “bulk” properties of (hyper)nuclear systems[16].

Figure 4 illustrates the influence of a growing contents of strangeness on the rms radii in ${}^{16+}O$. A considerable increase of the Λ rms radius in contrast to rather moderate changes of the total radius indicates that the loosely bound Λ 's in the systems with larger n_{Λ} tend to fill the surface region and thus produce a “halo” surrounding the core nucleons[11, 12].

The dependence of the binding energy per particle E/A on the number of hyperons for ${}^{16+}O$ and ${}^{40+}Ca$ is shown in Fig. 5. The behaviour of E/A results from a successive occupation of Λ -shells. There is apparently tighter binding when Λ 's fill deeply bound states, where the binding energy B_{Λ} is larger than E/A . After reaching the maximum for $n_{\Lambda} = 2$ (O) and $n_{\Lambda} = 8$ (Ca), the binding decreases as Λ particles start to fill higher orbits. The initial increase of the binding is a nice and clear demonstration of the absence of Pauli blocking between hyperons and nucleons, which overcomes the effect of weaker ΛN interaction compared with the NN one.

First studies of multi-strange objects have limited themselves to calculations of ΛN systems[11–13]. However, Schaffner *et al.*, [15] have recently pointed out that the considerations can not be restricted to nucleons and Λ particles only since Ξ^0 and Ξ^- are necessary constituents of the strange baryonic matter. This is because the process $\Lambda\Lambda \rightarrow \Xi N$ becomes energetically favourable due to the Pauli blocking of Λ 's in some multi- Λ hypernuclei. This greatly increases the possible amount of strangeness in the bound systems (stable against strong decay). The RMFT calculations predict baryonic (N, Λ, Ξ) systems with densities $\rho \approx (2-3)\rho_0$, $|S|/A \approx 1$ and $|Z|/A \ll 1$. These values resemble the predictions for droplets of strange quark matter “strangelets”[25]. However, the baryonic objects are more loosely bound ($E/A \approx -10 - -20$ MeV) and have shorter lifetimes. The reader can find an excellent review on the (N, Λ, Ξ) systems together with a detailed discussion of various baryon-baryon couplings in Ref. [15].

V. HYPERON-NUCLEUS SCATTERING

In the previous sections we have applied the RMFT to the description of hypernuclei and bound multi-strange systems. It is known from ordinary nuclei that the Dirac approach is a valuable tool to describe not only the nuclear shell model but the optical model, as well[6]. It is thus quite natural to extend our considerations to higher energies and try to apply Dirac phenomenology to Λ and Σ scattering off nuclei.

The optical potential for nucleons has been derived from the nucleon-nucleon t -matrix[55] or determined by a fit[56, 57] to nucleon-nucleus scattering data over a wide range of energies and nuclei. In this work we have chosen as a starting point the latter approach. We adopted the strengths and shapes for the real and imaginary parts of the nucleon-nucleus scalar and vector potentials from the global optical model of Cooper *et al.*[57]. For the real parts of the hyperon optical potential we assume the strengths scale like the mYY ($m=\sigma, \omega$ and $Y = \Lambda, \Sigma$) couplings.

The scaling factors α_{YV} used in our calculations are listed in Table IV. Following the ideology of Section 2 we determined the first set ($S1$) of the scaling factors α_{SY} , α_{VY} ($Y = \Lambda, \Sigma$) by searching for parametrization consistent with data on hyperon binding and simultaneously close to the $SU(3)$ values. For the sake of comparison we made calculations of the Y -nucleus scattering observables also for values of α_{YV} close to $1/3$. They are denoted as $S2$ in Table IV. The fits of the α_{YV} factors for real parts of the optical potentials were done using the geometry obtained from the global optical potential at 30 MeV[57].

We also include an anomalous coupling for the vector field with $f = -g$ and $f = +g$ for Λ and Σ , respectively. The tensor coupling term directly affects the strength of the hyperon spin-orbit interaction and therefore influences the analyzing power A_y . Since we limit our discussion to $N = Z$ targets we do not include the isovector ρ meson which can in principle contribute to the Σ optical potential[20].

The imaginary part of the optical potential arises predominantly from processes where the projectile interacts twice with the nucleus. Therefore, we use for the imaginary part of the optical potential the square of the factor for the real part. This is admittedly crude but should serve as a first estimate. In fact, the imaginary part will be increased by a substantial contribution from $\Sigma N \rightarrow \Lambda N$ conversion which was neglected here.

The α_{YV} values $S1$ from Table IV for potentials with optical model geometry and those from mean field calculations in Section 3 are close. It should be noted that this holds also for the other nuclei for which we obtained almost identical parametrization of the optical potential. This suggests that both the optical potential[57] and mean field model have very similar radial forms. Moreover, this serves as a check of our approach since we are not far from where the couplings in Section 2 have been determined.

Calculations for the differential cross sections and analyzing powers[30] were carried out using the modified code RUNT[58]. Here we present only the results for Λ and Σ hyperon scattering off ^{16}O at 200 MeV.

The comparison of the predictions of the parametrization $S1$ from Table IV for ^{16}O and different types of projectiles at 200 MeV is made in Fig. 6. The cross sections are qualitatively similar, small deviations are caused by the different contribution of the Coulomb interaction and tensor coupling for each type of hyperon. The predictions of A_y differ considerably between Σ and Λ scattering while the difference between Σ^+ and Σ^- is not so pronounced. This signals that A_y is affected far more by the tensor coupling than by the Coulomb interaction in this case. In order to confirm this observation we added in Fig. 6 the predictions for Λ with $f/g = 0$, as well. The $f/g = -1$ (value suggested by $SU(3)$) leads to a considerably

smaller values of analyzing power - A_y is almost zero for forward angles. This result suggests that measurements of A_y would give information on tensor coupling of the vector meson to a hyperon.

Figure 7 illustrates the sensitivity of scattering observables to the strength of the hyperon coupling. The calculations were done for $\alpha_{YV} \approx 1/3$ ($S2$) which is close to values frequently used in the literature[8, 11, 14, 21]. The angular dependence of the cross section is more or less similar to that of Fig. 6 - the differences for all hyperons start to appear at larger angles where the cross section is more than 2 orders of magnitude lower than for forward angles. The apparent bump in the cross sections for Σ^+ and corresponding sharp peak in A_y suggest a growing importance of the Coulomb interaction in this particular case. For Λ ($f/g = -1$) we obtained again very small A_y . For the other values of tensor coupling ($f/g = 0$ for Λ and $f/g = -1$ for Σ 's) the A_y differs significantly from the predictions of Fig. 6, even for small angles. The measurement of the analyzing power for Σ -nucleus scattering might thus bring some information about the strength of the YN interaction. On the other hand the small spin-orbit force due to $f = -g$ in the case of Λ -nucleus scattering means that we can not extract from A_y separate information about the scalar and vector potentials.

Calculations of Ref. [30] revealed that the results tend to be sensitive to the choice of input parameters. However they are similar enough for qualitative conclusions. The differential cross sections show the typical diffractive shape. The different parameter sets predict relatively similar angular dependence - the magnitudes can differ by more than a factor of two particularly at large angles. In the case of weak Σ^+ hyperon couplings, the Coulomb interaction might cause some interference effects in the scattering observables.

We observed strong sensitivity of A_y to the tensor coupling. The analyzing powers are small for the Λ ($f/g = -1$) but substantial for the Σ 's at energies around 200 MeV. Moreover, the predictions of A_y are quite sensitive to the strength of the Σ coupling. The measurement of A_y for Σ -nucleus scattering thus might become a source of information on YN interaction. However, it should be stressed again that $\Sigma N \rightarrow \Lambda N$ conversion was neglected here. Its contribution to the imaginary part of the optical potential could significantly affect the calculated analyzing powers for Σ . Therefore, this open problem deserves further study.

Currently there is no data on hyperon-nucleus scattering. However, this may change in the near future when an experiment at KEK may be able to provide data on Σ^+ -nucleus[31] scattering. The description of quasi-free hyperon production ($\gamma, K^+ \Lambda$), an approved proposal at CEBAF[59], also requires knowledge of the hyperon-nucleus optical potential.

The present calculations should be regarded as a first attempt to obtain a qualitative estimate of the hyperon-nucleus scattering observables. It is clear that much work, both theoretical and experimental, still remains in this field.

VI. NON-MESONIC DECAYS OF Λ HYPERNUCLEI

The weak decay of the Λ and other hyperons has been studied for almost 40 years as an example of a non-leptonic weak process. In the absence of strong interaction

the weak quark-quark interaction is assumed to be given by the standard model. However the modifications due to the strong interaction are not well understood. In nuclei, because of the Pauli suppression, the non-mesonic reaction $\Lambda N \rightarrow NN$ dominates[60] except for $A < 5$. It is this process that we will study in this section. Fortunately experimental information is available on both decay rates[61] and the asymmetry[62]. The discussion here follows that of Ramos *et al.*,[32, 33]

In the pion exchange description of the $\Lambda N \rightarrow NN$ the weak vertex Hamiltonian is given by:

$$H_\pi = ig_\omega \bar{\Psi}_N (1 + \lambda \gamma_5) \vec{\tau} \cdot \vec{\Phi}_\pi \Psi'_\Lambda, \quad (13)$$

where λ and g_ω are the empirically determined strengths. The nucleon and pion fields are given by $\bar{\Psi}_N$ and Φ_π respectively, while $\Psi'_\Lambda = \begin{pmatrix} 0 \\ \psi_\Lambda \end{pmatrix}$ is the Λ spurion field which is used to enforce the empirical $\Delta I = 1/2$ rule. The dynamics of the decay is contained in the relativistic two-body matrix element:

$$t_{ji} = g_{NN\pi} g_\omega \int d^4x d^4y \bar{\Psi}_{h_1, m_1, m_1}(x) (1 + \lambda \gamma_5) \vec{\tau}_i \Psi_{h_2, m_2, m_2}(y), \quad (14)$$

$$\times i \Delta_\pi(x - y) f(x - y) \bar{\Psi}_{h_2, m_2, m_2}(y) \gamma_5 \vec{\tau}_j \Psi_{h_1, m_1, m_1}(y), \quad (15)$$

where Ψ_{h_1, m_1, m_1} is the distorted outgoing nucleon wave function, $f(x - y)$ is the ΛN two-body correlation and $\Delta_\pi(x - y)$ is the combined pion propagator and vertex functions. The baryon wave functions, both bound state and scattering, are obtained from relativistic mean field approach.

We assume the weak coupling scheme for the hypernuclear wave function. Since a two baryon state is needed the core state is further decomposed using fractional parentage coefficients. The full details are in Ref. [33]. The final expression for the decay rate can then be written:

$$\Gamma_{nm} = \frac{1}{(2\pi)^3} \int d^3k_1 \int d^3k_2 \delta(M - E_f - E_i - E_2) \frac{M_N^2}{E_1 E_2 (2J_i + 1)} \sum |M_{fi}|^2, \quad (16)$$

where the sum runs over the initial polarizations and all the final state quantum numbers. The results are given in Table V for a few nuclei and different assumptions in the calculations.

The asymmetry can also be calculated (Table VI). It is given by

$$A = P_\nu \frac{3}{J + 1} \frac{\sum M_i \sigma(M_i) M_i}{\sum M_i \sigma(M_i)} \cos \chi = P_\nu A_\nu \cos \chi, \quad (17)$$

where χ is the angle between the proton direction and the polarization axis, $\sigma(M_i)$ is the intensity of protons along the quantization axis for a hypernucleus with spin projection M_i and P_ν is the polarization of the hypernucleus. It is necessary to calculate A_ν and determine P_ν from the experimental conditions[33].

From Tables V and VI we can see that while the asymmetry agrees within the large error bars the decay rates are significantly too low. On the other hand, it is interesting that the naive calculation without correlations (SRC) or form factors (FF) gives the decay rates larger than the experimental values. While adding kaon exchange has a significant effect the result is still way too small. Presumably it is necessary to also include the exchange of vector mesons ρ , ω and K^* . However, the problem is that the weak couplings of these mesons are not known and with so many additional parameters available one can fit everything while learning very little. It appears necessary to develop reliable models for the weak meson baryon couplings before more progress can be made.

VII. CONCLUSION

In this work we aimed to demonstrate that the relativistic mean field approach provides a natural description of the Y -nucleus systems. With reasonable (SU(3) inspired) values of the scalar, vector and tensor couplings it is possible to reproduce the experimental hypernuclear data: central potential, spin orbit interaction, single particle energies. The inclusion of the $\omega Y Y$ anomalous coupling into considerations has large consequences to the hypernuclear magnetic moments – for a Λ in the s state the originally predicted deviations between the relativistic predictions and Schmidt values disappear. The speculations about properties of multi-strange systems may seem too preliminary at present but, on the other hand, they have shown how challenging this study is. It can provide new information about baryon-baryon interaction and also distinguish between various theoretical approaches. Hyperon-nucleus scattering could become another important source of information about hyperon-nucleon interaction. The non-mesonic decay of a hypernuclei provides an example of where more work is required. Although there is agreement with the asymmetry the decay rates are seriously in error, perhaps indicating the need for additional meson exchanges.

Recent activities in experimental facilities with hypernuclear program – BNL, KEK, CEBAF, DUBNA – suggest the great perspectives of the hypernuclear studies in this decade.

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Table I. The parametrizations L [38] and SH [40] of the nucleonic sector. The masses are given in (MeV) and the coupling g_2 in (fm^{-1}).

	M_N	m_σ	m_ω	m_ρ	$g_{\sigma N}$	$g_{\omega N}$	$g_{\rho N}$	g_2	g_3
L	939.	520.0	783.	770.	10.47	13.80	8.07	0.0	0.0
SH	939.	526.065	783.	763.	10.444	12.945	8.766	-6.9099	-15.8337

Table II. The couplings used for hyperonic sector. The ratios of hyperon to nucleon couplings for $Y = \Lambda, \Sigma$ and Ξ are presented here. For $\alpha_{\pi Y}$ the values of Nijmegen models F[43], D[44], the values adopted from Cohen and Weber[18] and from Dover and Gal[3] are used for comparison.

	$\alpha_{\sigma Y} =$	$\alpha_{\omega Y} =$	$\alpha_{\rho Y} =$	$\alpha_{\pi Y} = f_{\omega Y}/g_{\omega Y}$
	$g_{\sigma Y}/g_{\sigma N}$	$g_{\omega Y}/g_{\omega N}$	$g_{\rho Y}/g_{\rho N}$	
$P1$	0.344	0.333	0.0	
Λ	$P2$ 0.462	0.500	0.0	0.0 -0.122[44] -0.541[43] -1.0 —
	$P3$ 0.621	0.667	0.0	
Σ	0.619	0.667	1.0	0.0 0.76 [3] — 1.0 1.417[44]
Ξ	0.375	0.333	1.0	0.0 -0.4 [44] -1.89[18,43] -2.0 -2.27 [18]

Table III. Hypernuclear magnetic moments for Λ in s state calculated using model L , parametrizations P1-P3 from Table II and $f/g = 0$. For $f/g = -1$ the predicted values (between -0.611 and -0.613 n.m.) are close to the Schmidt value.

	^{13}C	^{17}O	^{41}Ca	^{91}Zr	^{209}Pb
Schmidt			-0.613		
P1	-0.658	-0.649	-0.665	-0.669	-0.678
P2	-0.666	-0.656	-0.679	-0.683	-0.690
P3	-0.681	-0.669	-0.702	-0.710	-0.721

Table IV. The scaling factors α_{iY} ($i = S, V$ and $Y = \Lambda, \Sigma$) used in this work.

Particle	Line Type	α_V	i/g	α_S
Figure 1. ^{16}O 200 MeV - parametrization S1				
Λ	solid	0.667	0	0.621
Λ	dashed	0.667	-1	0.621
Σ^+	dash-dotted	0.667	1	0.616
Σ^-	dotted	0.667	1	0.616
Figure 2. ^{16}O 200 MeV - parametrization S2				
Λ	solid	0.333	0	0.347
Λ	dashed	0.333	-1	0.347
Σ^+	dash-dotted	0.333	1	0.345
Σ^-	dotted	0.333	1	0.345

Table V. Total decay rates for various hypernuclei (in units of free Γ_Λ)[33]. The experimental data are from Ref. [61].

	Γ_{nn}	$^5_\Lambda He(1/2^+)$	$^{11}_\Lambda B(5/2^+)$	$^{12}_\Lambda C(1^-)$
No correlations		1.00	1.38	1.62
SRC, FF		0.25	0.30	0.35
SRC, FF, distortions		0.16	0.25	0.29
π and K exchange		0.22	0.36	0.41
exp.		0.41 ± 0.14		1.14 ± 0.20

Table VI. Asymmetry $\mathcal{A}$ at $\chi = 0^\circ$ for polarized $^{12}_\Lambda C$ and $^{11}_\Lambda B$. The experimental data are from Ref. [62].

	$^{11}_\Lambda B(5/2^+)$	$^{12}_\Lambda C(1^-)$
	-0.108	-0.047
exp.	-0.20 ± 0.10	-0.01 ± 0.10

Figure Captions

1. Λ single particle energies for model SH (Table I), parametrization P3 (Table II) and $f_{\omega\Lambda}/g_{\omega\Lambda} = -1$ compared with experimental data[41].
2. Comparison of the Σ^- (upper part, a) and the Σ^+ (lower part, b) single particle levels for $\alpha_{\rho\Sigma} = 0$ and $\alpha_{\rho\Sigma} = 1$; $\alpha_{T\Sigma} = 1$.
3. Dependence of the hyperon single particle levels in ^{17}O ($Y = \Lambda, \Sigma^0, \Xi^0$) on the $\alpha_{TV} = f_{\omega V}/g_{\omega V}$ from Table II for a) $Y = \Lambda$, b) $Y = \Sigma^0$, c) $Y = \Xi^0$.
4. Dependence of the total (R_T) and Λ (R_Λ) rms radii in $^{16+u}_\Lambda O$ on the number of Λ hyperons (n_Λ) for model SH (Table I), parametrization P3 (Table II) and $f_{\omega\Lambda}/g_{\omega\Lambda} = 0, -1$.
5. Evolution of the binding energy per particle (E/A) in $^{16+u}_\Lambda O$ and $^{40+u}_\Lambda Ca$ for model SH (Table I), parametrization P3 (Table II) and $f_{\omega\Lambda}/g_{\omega\Lambda} = -1$.
6. The observables predicted for 200 MeV Λ (solid line for $f/g = 0$, long dash for $f/g = -1$), Σ^+ (dash-dotted line) and Σ^- (dotted line) elastically scattered from ^{16}O . The used parametrization is close to SU(3) and is given in Table IV.
7. The observables predicted for 200 MeV hyperons elastically scattered from ^{16}O for parametrization close to 1/3 from Table IV. Notation is that of Fig. 6.

Fig. 1

Fig. 3

Fig. 4

Fig. 2

Fig. 5

Fig. 6

Fig. 7