The Beauty of Spin

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Abstract. I review recent developments in theoretical spin physics. Topics include pion production in nucleon-nucleon collisions, the implications of heavy quark spin symmetry for heavy hadron molecules, the nucleon electric dipole form factors and ab initio calculations of the width of hadron resonances. A few spin physics high-lights from experiments at the COSY accelerator are also discussed.

1. Introduction and disclaimer
Spin is an intrinsic quantity of fundamental and composite particles. It is particularly fascinating since it is a quantum phenomenon - spin simply does not exist in classical physics. Furthermore, relativity combined with quantum mechanics leads to the spin-statistics theorem, which is yet another intriguing aspect of spin. Spin is also one of the finest experimental tools as it entails polarization. More precisely, experiments with polarized beams and/or targets allow to test the inner workings of the Standard Model, as will be discussed in a few selected examples below. Spin also entails a magnetic moment, and the comparison of its calculation and measurements has led to some of the most precise tests in physics, like e.g. for QED (see the recent paper [1] and references therein). Here, I can only address some topics related to spin and leave the discussion on the spin of the proton, the measurements and interpretation of single-spin asymmetries, the role of spin-dependent nuclear forces, or double polarization experiments at ELSA (Bonn), just to name a few, to other speakers at this symposium. In particular, an experimental overview on recent achievements in spin physics is given by Milner [2].

2. Theoretical spin physics at Jülich: a short overview
First, I would like to review some recent work on theoretical spin physics performed at Jülich. This is, of course, a very subjective choice but is intended to underline the importance of the spin degrees of freedom in hadron and nuclear physics.

Consider first the light quark sector of QCD, i.e. the two-flavor world of up and down quarks. Here, chiral perturbation theory (or related effective field theories) is the appropriate theoretical tool. Many processes have been analyzed in this framework, see e.g. Ref. [3]. The extension to systems with two or more nucleons requires some additional non-perturbative resummation, for a review see [4]. Recently, a combined analysis of the reactions $pn \rightarrow pp\pi^-$, $pp \rightarrow pn\pi^+$ and $pp \rightarrow d\pi^+$ at next-to-next-to-leading order (NNLO) in the chiral expansion became available [5]. In particular, the existing data were analyzed as a function of the three-nucleon low-energy constant (LEC) $D$ that parameterizes the leading $(NN)^2\pi$ contact operator.
Figure 1. Left: Results for the ratio of the Legendre coefficients in the threshold region for \( pp \to d\pi^+ \) as a function of the strength of the LEC \( D \). \( D = 3 \): solid, \( D = 0 \): dashed, \( D = -3 \): dot-dashed. Here, \( \eta \) is the pion momentum in units of the pion mass. Data: Ref. [9–12]. Right: Analyzing power \( A_y \) for \( pn \to pp(1S_0)\pi^- \). Data: Ref. [13–15].

It features prominently in a variety of low-energy processes like \( Nd \to Nd, NN \to NN\pi, NN \to dl\nu, \gamma d \to NN\pi \) or \( d\pi \to NN\gamma \). A precise determination of this LEC would therefore lead to stringent tests of the chiral QCD dynamics. In pion production reactions, one has to account for the scale \( \sqrt{M_\pi m} \simeq 340 \) MeV, with \( M_\pi \) the pion (nucleon) mass, in setting up the power counting, see [6] for a review. As can be seen from Fig. 1, the existing data are not precise enough to pin down \( D \) accurately, but seem to favor a positive value. Note that, not unexpected, the polarization observable \( A_y \) (right panel) is more sensitive to \( D \) than the threshold differential cross section, evaluated as \( d\sigma/d\Omega = A_0 + A_2 P_2(\cos \theta_\pi) \) (left panel) (for more details, see [7]). Therefore, at COSY an experiment has been proposed to determine \( D \) accurately from the measurement of \( \vec{p}\vec{n} \to \{pp\} \pi^- \) [8].

Now let me move to the heavy quark sector of \( c \) and \( b \) quarks, where heavy quark EFT is applicable. In the heavy quark limit \( m_Q \to \infty \), the QCD Lagrangian becomes independent of spin and flavor at leading order. Due to this symmetry, there are spin multiplets of both heavy mesons and heavy quarkonia, as e.g. the \( \{D,D^*\} \) and \( \{\eta_c,J/\psi\} \). The masses of the members within the same spin multiplet would be degenerate in the heavy quark limit. Recently, this symmetry was extended to possible heavy meson molecules [16], which were observed in recent years. Heavy meson molecules are bound states consisting of a heavy meson/heavy quarkonium and a light hadron, or two heavy mesons. It can be argued that the hyperfine splitting within a heavy quarkonium spin multiplet is untouched by their interactions with light mesons. As a result, a bound state of a heavy quarkonium and light hadrons, would have partner(s) whose components are the same light hadrons and the spin-multiplet partner(s) of the same heavy quarkonium. In particular, if the \( Y(4660) \) is a \( \psi' f_0(980) \) bound state [17], then there should exist an \( \eta_c' f_0(980) \) bound state, called \( Y_\eta' \), with its mass given as

\[
M_{Y_{\eta'}} = M_{Y(4660)} - (M_{\psi'} - M_{\eta_c'}) = 4616^{+5}_{-6} \text{ MeV.} \tag{1}
\]

The quantum numbers of such a state are \( J^P = 0^- \). The molecular picture of the \( Y(4660) \) describes naturally the invariant \( \pi^+\pi^- \) mass distribution measured in the reaction \( e^+e^- \to \gamma_{\text{ISR}} \pi^+\pi^- \psi' \) as shown in the left panel of Fig. 2. For other recent work on the molecular nature of the \( Y(4660) \), see Refs. [18,19]. Using heavy quark spin symmetry, one can then predict the analogous \( \pi^+\pi^- \) mass spectrum of the \( Y_\eta \) decay (right panel of Fig. 2). Its eventual measurement
would serve as an excellent testing ground of the molecular nature of both the \( Y\)\(_{4660}\) as well as its spin-partner, the \( Y\eta\). Exciting times are ahead of us in view of the many new data from BESIII at the BEPC, from the B-factories and, in the future, from PANDA at the HESR.

3. Experimental spin physics at Jülich: a few selected high-lights

Next, I would like to review some recent work on experimental spin physics performed at Jülich at the Cooler Synchrotron COSY. Again, this is a very subjective choice but is intended to show that COSY can be considered as “the spin machine”.

- COSY provides polarized proton and deuteron beams impinging on unpolarized and polarized targets. This allows to perform a variety of fundamental investigations. Besides the pion production experiments already mentioned, charge symmetry breaking in \( \vec{d} \, d \to \alpha \pi^0 \) (separation of s- and p-waves) will be studied with WASA-at-COSY. Furthermore, the reaction \( \vec{p} \, p \to pK\Lambda \) will be investigated at TOF with the aim of a determination of the singlet and triplet \( \Lambda p \) scattering lengths.

- The SPIN@COSY collaboration has performed many “spin-gymnastics” experiments. Of particular importance for a precision measurement of the \( \eta \) mass from the threshold cross section of the reaction \( pd \to ^3\text{He} \eta \) was the recent determination of the COSY beam momentum with an unprecedented accuracy, \( \Delta p/p < 10^{-4} \) at 3 GeV beam momentum, by depolarizing the beam through the use of an artificially induced spin resonance [21].

- The PAX collaboration works on methods to polarize antiprotons. It was recently suggested that this can be done in spin-flip reactions [22], although the underlying theoretical calculations were challenged early [23]. At COSY, the inverse reaction - namely the depolarization of protons through interactions with electrons - was measured and it was shown that the spin-flip mechanism advocated in [22] is not a viable tool to polarize antiprotons [24].

- In the context of the possible measurements of the proton and the deuteron electric dipole moment in storage rings (see the next section), the EDM collaboration performs a variety of tests at COSY, such a setting up and testing polarimeters to be used in EDM experiments [25].
4. The nucleon electric dipole moment

The neutron electric dipole moment (nEDM) is a sensitive probe of CP violation in the Standard Model and beyond. The current experimental limit \( d_n \leq 2.9 \cdot 10^{-26} \) cm is still orders of magnitude larger than the Standard Model prediction due to weak interactions. However, in QCD the breaking of the \( U(1)_A \) anomaly allows for strong CP violation, which is parameterized through the vacuum angle \( \theta_0 \). Therefore, an upper bound on \( d_n \) allows to constrain the magnitude of \( \theta_0 \). New and ongoing experiments with ultracold neutrons strive to improve these bounds even further, see e.g. [26] for a very recent review. On the theoretical side, first full lattice QCD calculations of the neutron and the proton electric dipole moment are becoming available [27–29]. These require a careful study of the quark mass dependence of the nEDM to connect to the physical light quark masses. In addition, there is a BNL proposal to measure the proton and the deuteron EDM in a storage ring [30] and there are also plans to build such type of machine in Jülich [31]. It is thus of paramount interest to improve the existing calculations of these fundamental quantities in the framework of chiral perturbation theory, as recently done in Ref. [32]. The calculation is based on \( U(3)_L \times U(3)_R \) covariant baryon chiral perturbation and includes all diagrams contributing at one loop order. To be specific, let me first give some basic definitions. The nucleon matrix element of the electromagnetic current in the presence of a magnetic field is given by

\[
\langle p'| J_{em} \nu | p \rangle = \bar{u}(p') \left[ \gamma^\nu F_1 (q^2) - \frac{i}{2m} \sigma^{\mu\nu} q_\mu F_2 (q^2) - \frac{1}{2m} \sigma^{\mu\nu} q_\mu \gamma_5 F_3 (q^2) + \ldots \right] u(p) \tag{2}
\]

with \( q_\mu = (p' - p)_\mu \). Here, \( F_1 \) and \( F_2 \) denote the P-, CP-conserving Dirac and Pauli form factors, \( m \) is the mass of the nucleon, and \( F_3 \) the P- and CP-violating electric dipole form factor. The ellipsis denotes the anapole form factor that we do not consider here. The electric dipole moment of the neutron/proton and the corresponding electric dipole radii follow as:

\[
d_{n,p} = \frac{F_{3,n,p}(0)}{2m} , \quad \langle r_{ed}^2 \rangle = \frac{6}{dq^2} \left. \frac{dF_3(q^2)}{dq^2} \right|_{q^2=0} . \tag{3}
\]

I summarize briefly the pertinent results of the study presented in Ref. [32]. First, a bound for the vacuum angle could be given, \( |\theta_0| \lesssim 2.5 \times 10^{-10} \). Second, the chiral expansion of the electric dipole radii takes the form (where \( \delta \) denotes a genuine small parameter)

\[
\langle r_{ed}^2 \rangle_n = -20.4 \left[ 1 - 0.67 + O(\delta^2) \right] \theta_0 \text{ cm fm}^2 , \\
\langle r_{ed}^2 \rangle_p = +20.9 \left[ 1 - 0.70 + O(\delta^2) \right] \theta_0 \text{ cm fm}^2 , \tag{4}
\]

where the large contribution of the NLO correction can be traced back to an enhancement of the pion loop effect due to an extra factor of \( \pi \), similar to what was observed in the analysis of the isospin-violating nucleon form factors [33]. Third, to compare results from (two-flavor) lattice QCD at unphysical quark masses with predictions from chiral perturbation theory, it is necessary to perform an extrapolation of the analytic results in the pion mass (see [32] for the details). In Fig. 3 we show the resulting pion (quark) mass dependence of the loop contribution to the neutron electric dipole moment in comparison to the available data points from two-flavor lattice QCD [28]. It is interesting to see that the complete one-loop calculation (solid line) reproduces the trend of the lattice data (the order of magnitude and the global sign) even without the unknown tree contribution, quite in contrast to the leading one-loop contributions (dot-dashed line). However, only below pion masses of the order of about 500 MeV the corrections are sufficiently small for a stable chiral extrapolation as indicated by the theoretical uncertainties also shown in Fig. 3. For other recent work on the electric dipole form factor of the nucleon, see [34,35].
It is also of interest to analyze the deuteron EDM, which has one- and two-body contributions. As shown in Ref. [36], the two-body corrections are generated from the induced P-wave in the deuteron wave function (polarization effects) as well as from pion-exchange currents. In particular, one finds a very different sensitivity to quark or chromomagnetic EDMs compared to the nucleon. For details, I refer to [36]. It certainly would be worthwhile to repeat this calculation in the framework of chiral nuclear EFT.

5. Ab initio calculation of hadron resonances

Arguably the most outstanding problem in QCD is to understand its spectrum. Questions in this context are: why does QCD mostly produce hadrons in terms of $\bar{q}q$ and $qqq$ states? What is the nature of the puzzling $X,Y,Z$ states and the related charm-strange mesons? Where are the exotica and the glueballs? Experimental facilities world-wide produce a cornucopia of high-precision data relevant to spectroscopy, often involving polarized beams and/or targets. In theory, lattice QCD is the premier tool to investigate the (low-lying) spectrum. In the light quark ($u,d,s$) sector, there has been tremendous progress with simulations at almost physical quark masses, see e.g. [37–40]. However, most hadrons are resonances - and therefore it is of utmost importance to be able to calculate their decay properties (partial and total width(s)).

I will concentrate here on the most significant baryon resonance, the $\Delta(1232)$. It is a well isolated resonance with its pole in the complex energy plane located at $(E_0,i\Gamma_0) = (1210,500)$ MeV. The $\Delta$ was first observed in pion-nucleon scattering and is known to dominate the photonuclear response. It is also amenable to a systematic EFT treatment, provided one counts the delta-nucleon mass splitting as a small parameter. The question how to calculate the width of resonance in a finite volume was in principle answered by Lüscher (and others) - in the vicinity of a narrow, well-separated resonance the two-particle energy levels exhibit a marked volume dependence in form of an avoided level crossing. This method was applied to synthetic finite volume data for pion-nucleon scattering in the $P_{33}$ partial wave in Ref. [41]. In fact, the avoided level crossing is washed out, but still a precise measurement of the volume dependence of the difference of the first two energy levels would allow to extract the delta-pion-nucleon coupling constant $g$ and thus the width of the $\Delta$. A different way of representing the same physics was discussed in Ref. [42] which in principle works for both narrow as well as broad resonances. This was demonstrated very nicely in a toy model by Morningstar [43] (see also Ref. [44]). Furthermore, the finite volume corrections for a decaying $\Delta$ have been worked out in [45] and successfully applied in combined fit to the data of the ETM collaboration for the pion mass dependence of the nucleon and the delta masses [46].
While these methods are working in principle, not sufficiently many lattice data for large enough volumes exist to cleanly extract the $\Delta$ width based on the energy level difference or the statistical method. However, for elastic two-body resonances, like the $\Delta$, a simpler method, originally proposed by Lüscher [47, 48] and Wiese [49], is based on computing the scattering phase shift in the infinite volume from the volume dependence of the energy levels of the lattice Hamiltonian. The delta resonances is practically a two-body state, and phenomenologically its scattering phase shift $\delta$ is very well described by the effective range formula

$$\frac{k^3}{W} \cot \delta(k) = \frac{6\pi}{g^2} \left( m^2 - W^2 \right),$$

where $k = |\vec{k}|$ is the constituents center-of-mass momentum, $W = \sqrt{k^2 + m_1^2} + \sqrt{k^2 + m_2^2}$ and $g$ is the coupling constant, which is related to the width of the resonance by

$$\Gamma = \frac{g^2 k^3}{6\pi m^2}.$$  

The phase shift $\delta$ passes through $\pi/2$ at the physical mass, i.e. $W^2 = m^2$. In the case of noninteracting particles the possible energy levels in a periodic box of length $L$ are given by

$$W = \sqrt{k^2 + m_1^2} + \sqrt{k^2 + m_2^2},$$

where $k = 2\pi |\vec{n}|/L$, $\vec{n}$ being a vector with components $n_i \in \mathbb{N}$. In the interacting case, the energy levels are still given by (7), but now $k$ is the solution of [42, 47, 48]

$$\delta(k) = \arctan \left\{ \frac{\pi^{3/2} \sqrt{q}}{Z_{00}(1; q^2)} \right\} \mod \pi, \quad q = \frac{kL}{2\pi},$$

where $Z_{00}$ is a generalized zeta function,

$$Z_{00}(1; q^2) = \frac{1}{\sqrt{4\pi}} \sum_{\vec{n}\in\mathbb{Z}^3} \frac{1}{\vec{n}^2 - q^2}.$$  

That is to say, each energy value $W$, computed on the periodic lattice at some fixed values of $m_1, m_2$, gives rise to a certain momentum $k$. The scattering phase at this momentum and pion mass is given by Eq. (8). Fitting $\delta$ to the effective range formula Eq. (5) then allows us to estimate the mass and width of the actual resonance. Eq. (8) holds for vanishing total momenta. For nonvanishing momenta different zeta functions or combinations of zeta functions apply [50].

The $\pi N$ phase shift of the $\Delta$ channel is more difficult to compute than the corresponding $\pi\pi I = J = 1$ phase that contains the $\rho$ (see e.g. Refs. [51–54]) because the phase space is much smaller at comparable pion masses. Nonetheless, together with the QCDSF collaboration, we are presently investigating the $I = 3/2$ P-wave in $\pi N$ scattering, using a $N_f = 2$ clover action that consists of the plaquette gluon action together with nonperturbatively $O(a)$ improved Wilson (clover) fermions. In Fig. 4 we show some first results, combining calculations on $32^3 \times 64$ and $40^3 \times 64$ lattices at $\beta = 5.29, \kappa = 0.13632, a = 0.075$ fm and on $40^3 \times 64$ and $48^3 \times 64$ lattices at $\beta = 5.29, \kappa = 0.1364, a = 0.06$ fm. So far the errors are relatively large, which calls for higher statistics. To trace the phase shift over a sufficiently large range of values around $\delta = \pi/2$ with high precision, calculations need to be formed on larger lattices and with nonvanishing total momenta. Such simulations are presently being performed. Still, as indicated by the solid line that is generated with the physical mass and width of the $\Delta$, we definitely are on the right track [55].
6. Summary & outlook
Spin physics is and will be an exciting field of research. Much progress has been made through developments in accelerator, experimental and theoretical physics. Here, I have focused on some recent developments in the sector of non-perturbative QCD, stressing the fruitful interplay between theory and experiment. On the theoretical side, much progress has been made in the use and application of effective field theories and/or lattice simulations. While lattice simulations are quickly approaching the limit of physical pion masses, chiral extrapolations and finite volume EFTs are nevertheless required to fully access the underlying physics, as best exemplified for the case of unstable particles on the lattice. Similarly, the rich spectrum in the charm sector measured at colliders and B-factories requires a thorough investigation of possible mechanisms to generate hadrons, such as the formation of hadronic molecules discussed in Sec. 2. Finally, precision measurements combined with accurate theoretical tools allow for many fine test of the chiral QCD dynamics, as discussed here for the determination of the low-energy constant $D$, that appears in a cornucopia of low-energy processes. The exciting possibility of measuring the electric dipole moment of the proton and light nuclei should lead to refined theoretical investigations, especially using chiral nuclear effective field theory. Experiments involving polarization will continue to deliver accurate data that will challenge our understanding of the Standard Model and might give possible signatures of physics beyond it. In this context, it would be very important to understand the mechanism behind the polarization of antiprotons, which, if implemented within the HESR complex at FAIR, would open exciting new possibilities to explore the physics of charm and light quarks in hadronic and nuclear systems.

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