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Resonances, Polarizations and Symmetries in Λ_b Decays [*]

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Abstract

Relativistic polarization offers very interesting opportunities to test fundamental symmetries like Time-Reversal (TR). Starting from simple scattering-matrix relations, supplemented by QCD-inspired models describing heavy quark decays, it is shown that the polarization-vectors of resonances coming from Λ_b decays exhibit very clear possibilities of TR violation.

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1 Introduction

The present letter is a complement to two preceding ones [1, 2] devoted to the study of Λ_b decays into $\Lambda V(1^-)$ where tests of both direct CP violation and Time-Reversal were developed. Our method is built on the procedure of cascade decays, for instance $\Lambda_b \to \Lambda J/\psi$, $\Lambda \to p\pi^-$ and $J/\psi \to \mu^+\mu^-$, procedure which requires a precise determination of the quantum properties of the intermediate resonances. In this kind of decays, resonance polarizations play crucial role for determining the angular and energy distributions of the final decay products. Moreover, the polarization-vectors have specific features when physical systems are transformed by discrete symmetries, like Parity and Time-Reversal. Since the spin of any particle is Parity-even and T-odd, studying its polarization provides interesting tools to test both of these two symmetries, especially in weak decay processes where parity is known to be violated. It is worth noticing that checking TR in a given decay, like $A \to a_1 + a_2$, or in its charge conjugate mode, $\bar{A} \to \bar{a}_1 + \bar{a}_2$, is not necessarily related to the conservation or non-conservation of the CP symmetry in these two channels. A direct check of TR is proposed, without referring to the CPT theorem.

According to the important role of the spin in particle physics and to the existence of different approaches to deal with the polarization problem, we suggest to clarify these notions in the present letter. We emphasize the relativistic aspect of spin in order to justify the intensive use of the kinematic formalism by Jacob-Wick-Jackson [4, 5] to previous analysis of cascade decays [1, 2, 3]. These applications are non-trivial, since, as we shall see, they involve Lorentz boosts. On the dynamical side, the Heavy Quark Effective Theory [6, 7, 8] (HQET) formalism is used to evaluate the hadronic form factors involved in Λ_b-decay. Weak transitions including heavy quarks can be safely described when the mass of a heavy quark is large enough compared to the QCD scale, Λ_{QCD}. Properties such as flavor and spin symmetries can be exploited in such a way that corrections of the order of $1/m_Q$ are systematically calculated within an effective field theory. Then, the hadronic amplitude of the weak decay is investigated by means of the effective Hamiltonian, $\Delta B = 1$, where the Operator Product Expansion formalism separates the soft and hard regimes. All results about transition form factors as well as hadronic matrix elements are given in Ref. [2].

Our analysis is oriented towards some special aspects of relativistic kinematics in the cascade decays of resonances. After briefly summarizing, in Section 2, the basic ingredients of relativistic spin, we introduce the particular frames in which polarization vectors of resonances and their physical properties are studied: mainly their transformation under Parity and Time-Reversal. In Section 3, computations of the polarization vectors are reviewed, stressing on the important role of the Λ_b density-matrix. In Section 4 detailed numerical results are displayed according to various input parameters. Emphasis is also put on the physical observable, the normal component of the polarization vector, which could clearly show possible violation of TR. Lastly a short conclusion is drawn in sect. 5.

2 Relativistic form of the polarization

2.1 Basics of the Relativistic Spin

The spin vector operator of a massive particle, $\vec{s} \equiv (s_1, s_2, s_3)$, whose components verify the standard commutation relations, $[s_i, s_j] = i\epsilon_{ijk}s_k$, is defined in the rest frame of the particle. A covariant extension of this operator is the Pauli-Lubanski four-vector, (S^α), defined in such a way that, in the particle rest frame, $S' \equiv (S'^0 = 0, \vec{S'} = \vec{s})$. This quadrivector fulfills the covariant constraint [9], $p_\mu S^\mu = 0$ where p_μ is the particle 4-momentum and, in the particle rest-frame, $p_\mu \equiv (m, \vec{0})$.

So, (S^α) behaves as an ordinary quadrivector when it is transformed from a Lorentz frame to another one. Setting $S = (S^0, \vec{S})$ in a given frame (R) which moves with the velocity $\vec{\beta}$ with respect to the particle, the components of S transform in the following way [10]:

$$S^0 = \gamma\vec{\beta}\cdot\vec{s}, \; \vec{S} = \vec{s} + \frac{\gamma^2}{\gamma+1}(\vec{\beta}\cdot\vec{s})\vec{\beta}, \qquad \text{with } \gamma = 1/\sqrt{1-\beta^2} \; .$$

Thus, knowing the spin vector of a particle in its rest-frame allows us to derive the Pauli-Lubanski quadrivector in any other frame.

These basic notions can be easily extended to the particle polarization vector, defined as the mean value of its spin vector in its own rest-frame, $\vec{\mathcal{P}} = \langle\vec{s}\rangle$. In this connection we observe that, according to Ehrenfest's theorem concerning the expectation value of a quantum observable, the polarization vector $\vec{\mathcal{P}}$ behaves as a classical quantity and its components are transformed as indicated by the formulas above, when performing Lorentz transformation between two different frames. It is worth noticing, too, that the time-component of the polarization vector of a massive spinning particle vanishes in its rest-frame, while it is generally different from zero when this particle is moving.

All these considerations can be applied to the resonances coming from Λ_b decays into $\Lambda V(1^-)$, the beauty baryons and especially the Λ_b are expected to be produced very copiously in proton-proton collisions with the next LHC machine.

2.2 Choice of particular frames

We consider decays of the type $\Lambda_b \to R_{(1)}R_{(2)}$, where the $R_{(i)}$ $(i = 1,\, 2)$ are resonances. The tests we propose are suitably performed by introducing, in the Λ_b rest-frame[5], a specific frame for each resonance $R_{(i)}$, according to Jackson's method:

$$\vec{e}_L = \frac{\vec{p}}{p} \; , \quad \vec{e}_T = \frac{\vec{e}_Z \times \vec{e}_L}{|\vec{e}_Z \times \vec{e}_L|} \; , \quad \vec{e}_N = \vec{e}_L \times \vec{e}_T \; , \tag{1}$$

where $\vec{e_Z}$ is parallel to $\vec{n}$; $\vec{n}$ being initially defined as the normal unit-vector to the Λ_b production plane. Then, each polarization-vector can be expressed as

$$\vec{\mathcal{P}} = P_L\vec{e}_L + P_N\vec{e}_N + P_T\vec{e}_T \; , \tag{2}$$

where P_L, P_N and P_T are respectively the longitudinal, normal and transverse components of $\vec{\mathcal{P}}$. It is useful to notice that the basis vectors $\vec{e_L}, \vec{e_T}$ and $\vec{e_N}$ have the following properties

[5]The Λ_b rest frame used in our analysis is given in Figure 1.

according to parity and TR: P-odd,T-odd; P-odd,T-odd and P-even, T-even respectively, while the polarization-vector $\vec{\mathcal{P}}$ is P-even and T-odd. So, any component of $\vec{\mathcal{P}}$ defined by the scalar product $P_j = \vec{\mathcal{P}} \cdot \vec{e}_j$ with $j = L, N, T$ gets transformed as:

$$P_L = \mathrm{P-odd, T-even}, \;\; P_T = \mathrm{P-odd, T-even} \text{ and } P_N = \mathrm{P-even, T-odd}.$$

Note that the longitudinal axis defined by $\vec{e_L}$ is taken as the quantization axis; $\vec{e_N}$ and $\vec{e_T}$ are identified respectively to x and y axis.

In order to measure $\vec{\mathcal{P}}^{(i)}$ for each resonance, $R_{(i)}$, and to determine its transformation by symmetry operations like Parity and Time-Reversal, a thorough examination of the $R_{(i)}$ decay products by sophisticated methods (Byers, Dalitz) [11] must be done in the resonance rest-frame itself. We shall not perform numerical simulations of such decays in the present paper. However we observe that, in the cascade decays we are concerned with, from the experimental point of view, the above mentioned basis vectors are more suitably defined in the Λ_b rest frame than in the one of $R_{(i)}$. Therefore we have to perform a Lorentz transformation. In particular, we shall apply a rigorous method, using the relativistic spin, in order to understand the modification of $\vec{\mathcal{P}}^{(i)}$ and, especially, its normal component P_N, when considering two different rest-frames: the Λ_b one and the $R_{(i)}$ resonance one, which are related to each other by a simple Lorentz boost. Thus, we are led to study the spin of any particle and its polarization in their relativistic form. In what follows, we will drop the index (i) which designates any resonance (Λ or J/ψ) and we will assign a prime to the physical quantities defined in the $R_{(i)}$ rest-frame.

2.3 Different rest-frames

According to the considerations of Subsect. 2.1, we denote by $\mathcal{P}' \equiv (0, \vec{P'})$ the polarization four-vector of the Λ hyperon in its own rest-frame. On the other hand, we define, in the rest-frame of Λ_b, the four-vectors $e_N \equiv (0, \vec{e}_N)$ and $e_T \equiv (0, \vec{e}_T)$, where the unit vectors $\vec{e}_N$ and $\vec{e}_T$ have been previously defined in the same frame. Since these unit vectors are orthogonal to relative momentum of Λ with respect to Λ_b, they are left unchanged in the boost from the Λ_b rest frame to the one of Λ. Then we have, owing to Lorentz invariance of the scalar product of two four-vectors,

$$P_N = -\mathcal{P} \cdot e_N = -\mathcal{P}' \cdot e'_N = P'_N, \tag{3}$$

where we have denoted by $\mathcal{P}$ the Λ polarization four-vector in the Λ_b rest frame. Then we conclude that the normal component, P_N, of the polarization vector is the same in the two different frames considered. An analogous reasoning leads us to stating that also the transverse component, P_T, shares the same interesting physical property. In particular, as regards P_N, this fact allows us to cross-check TR symmetry either in the Λ_b rest-frame or in any resonance rest-frame coming from Λ_b decays.

3 Polarizations of the final resonances

Basic principles of Quantum Mechanics allow us to deduce the spin density matrix of the final ΛV system, which is an essential parameter to compute the polarization-vector of each resonance, $\rho^f = \mathcal{T}^\dagger \rho^{\Lambda_b} \mathcal{T}$. $\mathcal{T}$ is the transition-matrix (the $\mathcal{S}$-matrix being defined by $\mathcal{S} = 1 + i\mathcal{T}$) whose elements are explicitly given in Ref. [2]. The normalization of the matrix ρ^f is obtained by

$$Tr(\rho^f) = \frac{d\sigma}{d\Omega} = NW(\theta, \phi), \tag{4}$$

where Tr is the trace operator and N is a normalization constant. Consequently, the polarization-vector of any resonance $R_{(i)}$ ($R_1 = \Lambda,\ R_2 = V$) is defined by

$$\vec{\mathcal{P}}_i = \langle \vec{S}_i \rangle = \frac{Tr\left(\rho_i^f \vec{S}_i\right)}{Tr\left(\rho_i^f\right)} \ , \tag{5}$$

where ρ_i^f is the spin density-matrix of the resonance, $R_{(i)}$, deduced from ρ^f. As the final state is a composite system made out of two particles with different spin ($s_1 = 1/2, s_2 = 1$), each ρ_i^f will be obtained from the general expression of ρ^f by summing up over the degrees of freedom of the other resonance. Thanks to this method, we can obtain ρ^Λ and ρ^V. See Ref. [2] for all analytical results.

The previous relation allows us to write down the formal expression of any $\vec{\mathcal{P}}_i$ by expanding the trace operator over the different spin states[6] and one gets,

$$\vec{\mathcal{P}}_i\, W(\theta, \phi) = N\ \Sigma_\lambda \Big(\langle \theta, \phi, \lambda | \rho_i^f \vec{S} | \theta, \phi, \lambda \rangle \Big) \ , \tag{6}$$

$W(\theta, \phi)$ being defined by Eq. (4).

The matrix elements given in the right-hand side of Eq. (5) can be explicitly calculated [5] and the three components of $\vec{\mathcal{P}}^\Lambda$ get the following expressions:

$$\begin{aligned} P_x^\Lambda\ W(\theta,\phi) &\propto 2\Re e(\langle \theta, \phi, 1/2 | \rho^\Lambda | \theta, \phi, -1/2 \rangle) \ , \\ P_y^\Lambda\ W(\theta,\phi) &\propto -2\Im m(\langle \theta, \phi, 1/2 | \rho^\Lambda | \theta, \phi, -1/2 \rangle) \ , \\ P_z^\Lambda\ W(\theta,\phi) &\propto \bar{\omega}(+1/2) - \bar{\omega}(-1/2) \ , \end{aligned} \tag{7}$$

where the $\bar{\omega}(\pm)$ are defined in [2]. The vector meson has spin $S_{(2)} = 1$ and therefore three helicity states. Based on

$$\vec{\mathcal{P}}^V\ W(\theta, \phi) = N\ \Sigma_{\lambda_2} \Big(\Sigma_{\lambda_1} \langle \theta, \phi, \lambda_1, \lambda_2 | \rho^f \vec{S} | \theta, \phi, \lambda_1, \lambda_2 \rangle \Big) \ , \tag{8}$$

the components of $\vec{\mathcal{P}}^V$ are obtained in the same manner, although more tedious. One has

$$\begin{aligned} P_x^V\ W(\theta,\phi) &\propto \sqrt{2}\Re e\big((\langle 0 | \rho^V | -1 \rangle) + (\langle 1 | \rho^V | 0 \rangle) \big) \ , \\ P_y^V\ W(\theta,\phi) &\propto \sqrt{2}\Im m\big((\langle 0 | \rho^V | -1 \rangle) + (\langle 1 | \rho^V | 0 \rangle) \big) \ , \\ P_z^V\ W(\theta,\phi) &\propto (\langle 1 | \rho^V | 1 \rangle) - (\langle -1 | \rho^V | -1 \rangle) \ . \end{aligned} \tag{9}$$

[6] In order to perform these calculations, some simple and fundamental relations are used, whatever the spin is:

$$S_x|\lambda\rangle = (|\lambda+1\rangle + |\lambda-1\rangle)/\sqrt{2}\ , S_y|\lambda\rangle = i(-|\lambda+1\rangle + |\lambda-1\rangle)/\sqrt{2}\ , S_z|\lambda\rangle = \lambda|\lambda\rangle \ .$$

(Letters indicating other physical parameters are dropped for simplicity). Technical details of the computation of both $\vec{\mathcal{P}}^\Lambda$ and $\vec{\mathcal{P}}^V$ are given in Ref. [2].

As it was expected, the helicity value $\lambda_V = 0$ does not contribute to the longitudinal polarization of the vector-meson. We also underline the importance of the initial Λ_b polarization, $\mathcal{P}^{\Lambda_b}$, as well as the non-diagonal matrix element, $\rho_{+-}^{\Lambda_b}$, in the analytical expression of the components of $\vec{\mathcal{P}}^V$.

4 Numerical Results and related problems

• In a first step, values for all input parameters are taken from [2]. The initial Λ_b polarization and Λ_b polarization density matrix element used in our numerical computations are $\mathcal{P}^{\Lambda_b} = 100\%$ and $\Re e(\rho_{+-}^{\Lambda_b}) = \Im m(\rho_{+-}^{\Lambda_b}) = \sqrt{2}/2$, respectively. The corresponding spectra of the three components of $\vec{\mathcal{P}}^\Lambda$ and $\vec{\mathcal{P}}^V$ are shown in Figures 2 and 3. Some comments on longitudinal, transverse and normal components of these polarization vectors are in order: (i) the longitudinal components, $P_L = P_z$, of both resonances are asymmetric because of parity violation in weak Λ_b decays; (ii) the spectra of the transverse components, $P_T = P_y$, are quite symmetric, their asymmetries being $\approx 1.0\%$; (iii) the normal components, $P_N = P_x$, are clearly asymmetric. Their asymmetry values are respectively 23% and -54% for Λ and J/ψ.

• In a second step, attempts to understand correlations between the Λ_b initial polarization and the physical properties of its decay products are made. All results are obtained with Monte-Carlo simulations by varying independently $\mathcal{P}^{\Lambda_b}$ and $\rho_{+-}^{\Lambda_b}$. Non-diagonal matrix elements being generally unknown, we set $\Re e(\rho_{+-}^{\Lambda_b}) = \Im m(\rho_{+-}^{\Lambda_b}) = 0$ and we let $\mathcal{P}^{\Lambda_b}$ vary between 100% and 0%. The resulting spectra of the normal components, P_N^Λ and $P_N^{J/\psi}$ are usually sharp, while the transverse components, both P_T^Λ and $P_T^{J/\psi}$, are always equal to zero. These two physical properties can be explained as direct consequences of $\rho_{+-}^{\Lambda_b} = 0$. In Table 1, mean values and asymmetry parameters of the normal component spectra for Λ and J/ψ are listed. Interesting remarks can be drawn: (i) for $\mathcal{P}^{\Lambda_b} \neq 0$, the normal components are largely dominating and their asymmetries are nearly equal to minus one; (ii) for $\mathcal{P}^{\Lambda_b} = 0$, the J/ψ normal component is still dominating (≈ -0.8) while P_N^Λ is equal to zero, the Λ polarization being completely longitudinal ($P_L^\Lambda = -100\%$).

• In order to understand the role of the Λ azimuthal distribution and its effects on the resonance polarization-vectors, a comparison between two series of P_N spectra is performed. One series is obtained with $\Re e(\rho_{+-}^{\Lambda_b}) = \Im m(\rho_{+-}^{\Lambda_b}) = 0$, while the other one is obtained with the standard values, $\Re e(\rho_{+-}^{\Lambda_b}) = \Im m(\rho_{+-}^{\Lambda_b}) = \sqrt{2}/2$. In Figures 4 and 5, the spectra of P_N and P_T for Λ and J/ψ are respectively plotted. One notes that the spectra belonging to $\Re e(\rho_{+-}^{\Lambda_b}) = \Im m(\rho_{+-}^{\Lambda_b}) \neq 0$ are much broader than those belonging to $\rho_{+-}^{\Lambda_b} = 0$. Moreover, the absolute values of the asymmetry parameters decrease for both P_N^Λ and $P_N^{J/\psi}$ when $\rho_{+-}^{\Lambda_b} \neq 0$, while the transverse components remain symmetric with mean values around 0. Whatever the elements of the Λ_b spin density matrix are, this exhaustive study indicates that the normal components of the polarization-vectors, which are T-odd observables, must be taken as serious candidates to cross-check TR symmetry.

4.1 Final State Interactions

In the above calculations, no mention has been made about Final State Interactions, FSI, which play an important role in all hadronic interactions. As it is known, when hadrons are produced in any process, strong interactions among them could distort their wave-functions by generating both additional phase-shifts and absorptive effects [12]. This modifies physical observables like decay widths, asymmetries and, in our particular case, T-odd observables. So, fundamental questions may arise:

(i) If T-odd parameters are experimentally observed, do these effects originate from true physical dynamics which violate TR symmetry, or are they consequences of FSI which mimic TR violation? (ii) The above question can also be reversed: if, in some process where expected T-odd effects are not seen, does this situation show effects of FSI which inhibit T-odd parameters?

Recently, in the framework of deep inelastic scattering experiments, D. Sivers [13] emphasized the difference between the antiunitary TR operator and the "artificial" or "naive" TR, which is unitary. As far as we are concerned, we will try to clarify this issue and bring some insights about the particular promising channel $\Lambda_b \to \Lambda J/\psi$ which branching ratio is $\approx 10^{-4}$. A simple quantum-mechanical approach has been given long time ago by DeRujula et al. [14] and it has been developed again by the authors of reference [15] after the discovery of T-violation by CPLEAR experiment.

Defining the S-matrix as $S = 1 + iT$, and $|i\rangle$ and $|f\rangle$ being respectively the initial and final states, the unitarity condition, $SS^\dagger = 1$, leads to: $T^*_{fi} = T_{if} - iA_{if}$ where the new term $A_{if} = \Sigma_n T_{in} T^*_{fn}$ indicates the transition amplitude from the initial state $|i\rangle$ to the final state $|f\rangle$ by including scattering on the on-shell intermediate states $|n\rangle$. This term is called the absorptive part of the transition amplitude and it is the main contribution to the FSI.

Let $\bar{i}$ and $\bar{f}$ be respectively the initial and final states with reversed momenta and spins, then the expression $|T_{if}|^2 - |T_{\bar{i}\bar{f}}|^2$ indicates the probability of a T-odd effect, while the expression $|T_{fi}|^2 - |T_{\bar{i}\bar{f}}|^2$ corresponds to the standard TR violation probability. Using the above relation which defines A_{if}, a simple relation among the two preceding probabilities can be deduced:

$$|T_{if}|^2 - |T_{\bar{i}\bar{f}}|^2 = (|T_{fi}|^2 - |T_{\bar{i}\bar{f}}|^2) - 2\Im m(A_{if} T^*_{if}) - |A_{if}|^2 \ .$$

If TR is an exact symmetry, then $|T_{fi}| = |T_{\bar{i}\bar{f}}|$ and the term in parenthesis on the right-hand side vanishes; instead the other terms generate T-odd effects, even if TR is an exact symmetry.

We turn now to the experimental case, $\Lambda_b \to \Lambda J/\psi$ followed by the cascade decays $\Lambda \to p\pi^-$, $J/\psi \to \mu^+\mu^-$. Examining thoroughly this channel, we expect that FSI occur at the decay vertex $\Lambda_b \to \Lambda J/\psi$ where, at the partonic level, (virtual) gluons can be exchanged and absorbed by the different quarks entering the Λ_b decay mechanism. After hadronization, and owing to the long Λ life-time, the hyperon Λ will decay far away from the J/ψ and their decay products, p, π^- and μ^+, μ^- respectively, will interact only electromagnetically, which is an insignificant effect because of the long distance separating the two mother resonances.

Understanding the FSI among the Λ and J/ψ requires rigorous analysis of experimental data which must be compared to our phenomenological model; and we can assert that it is the most realistic solution to this crucial problem.

5 Conclusion

Complete calculations based both on the helicity formalism (kinematics) and on the OPE techniques supplemented by HQET (dynamics) have been performed in a rigorous way for a precise determination of the physical properties of the $\Lambda_b \to \Lambda V(J^P = 1^-)$ decays. Resonances Λ and $V(J^P = 1^-)$ being polarized, it is shown that the normal components of their polarization-vectors are T-odd observables. Furthermore, these components have large asymmetries and they are Lorentz-invariant. An exhaustive study of $P_N^{\Lambda, J/\psi}$ has been performed according to the Λ_b polarization density matrix. Thanks to our analysis, it is confirmed that these observables are truly serious candidates to cross-check Time-Reversal symmetry, and we hope to detect these effects with the forthcoming LHC machine.

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$\mathcal{P}^{\Lambda_b}$	P_N^{Λ}	As^{Λ}	$P_N^{J/\psi}$	$As^{J/\psi}$
100%	-0.98	-1.0	-0.88	-0.95
75%	-0.97	-1.0	-0.89	-1.0
50%	-0.96	-1.0	-0.87	-1.0
25%	-0.88	-1.0	-0.85	-1.0
10%	-0.61	-1.0	-0.83	-1.0
0%	0.0	0.0	-0.81	-1.0

Table 1: Mean values, $P_N^{\Lambda,J/\psi}$, and asymmetries, $As^{\Lambda,J/\psi}$, of the polarization-vector normal components of Λ and J/ψ, respectively. Results are given as functions of the initial Λ_b polarization varying from 100% to 0% and for the $\Lambda_b \to \Lambda J/\psi$ decay channel.

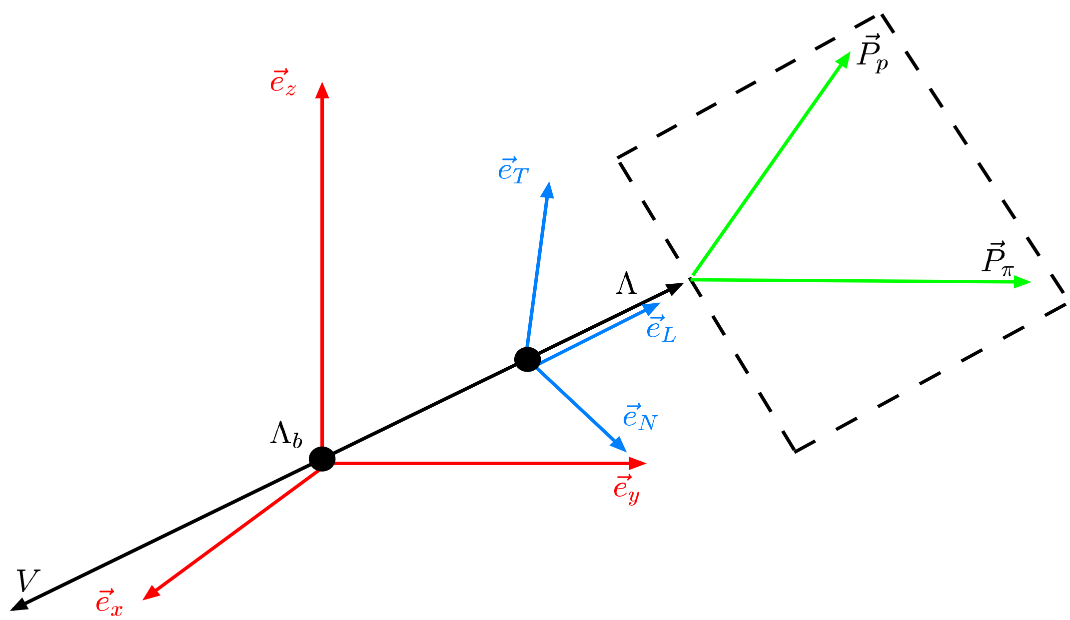

Figure 1: The $\vec{e}_x, \vec{e}_y, \vec{e}_z$ as well as the $\vec{e}_T, \vec{e}_N, \vec{e}_L$ frames in the Λ_b rest-frame.

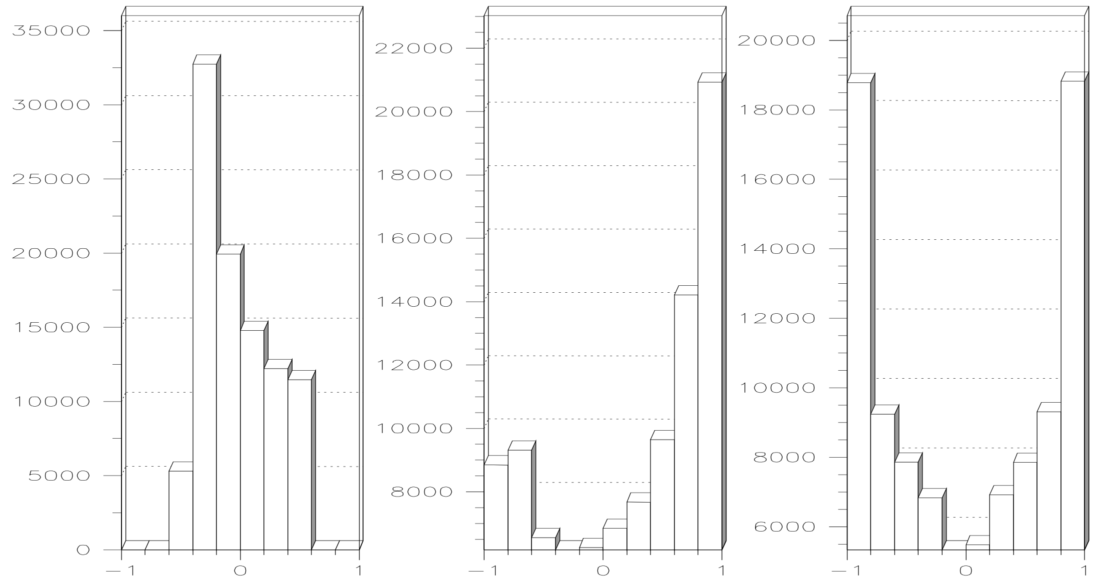

Figure 2: Spectra of the Λ polarization-vector components: (from left to right) P_L, P_N, P_T, respectively in case of $\mathcal{P}^{\Lambda_b} = 100\%$ and $\Re e(\rho_{+-}^{\Lambda_b}) = \Im m(\rho_{+-}^{\Lambda_b}) = \sqrt{2}/2$.

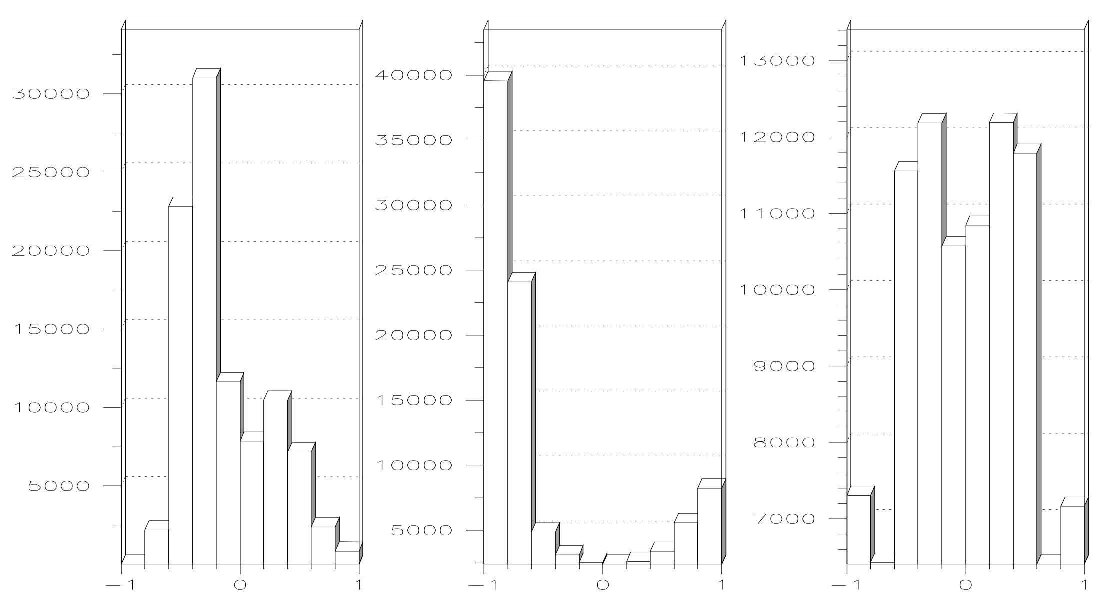

Figure 3: Spectra of the J/ψ polarization-vector components: (from left to right) P_L, P_N, P_T, respectively in case of $\mathcal{P}^{\Lambda_b} = 100\%$ and $\Re e(\rho_{+-}^{\Lambda_b}) = \Im m(\rho_{+-}^{\Lambda_b}) = \sqrt{2}/2$.

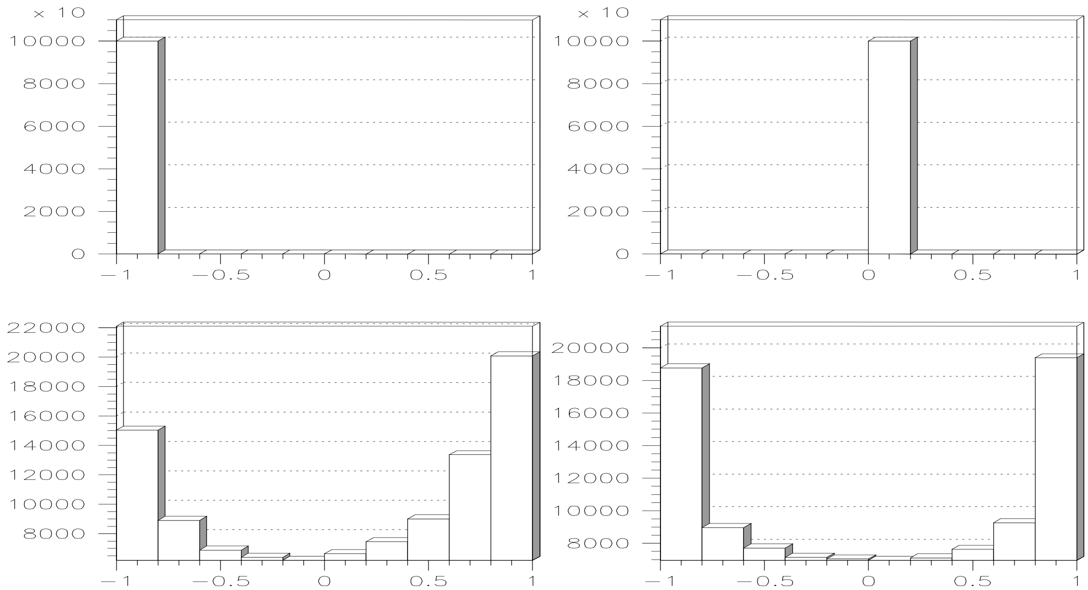

Figure 4: Spectra of P_N^Λ and P_T^Λ with $\mathcal{P}^{\Lambda_b} = 50\%$. Upper histograms correspond to the case of $\Re e(\rho_{+-}^{\Lambda_b}) = \Im m(\rho_{+-}^{\Lambda_b}) = 0$, while lower histograms correspond to $\sqrt{2}/2$.

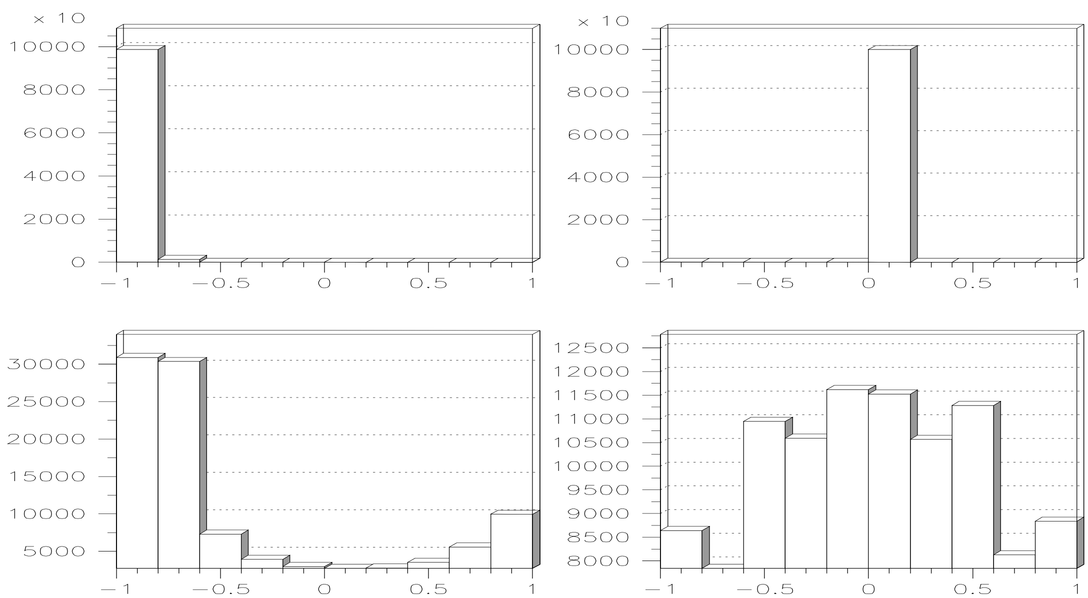

Figure 5: Spectra of $P_N^{J/\psi}$ and $P_T^{J/\psi}$ with $\mathcal{P}^{\Lambda_b} = 50\%$. Upper histograms correspond to the case of $\Re e(\rho_{+-}^{\Lambda_b}) = \Im m(\rho_{+-}^{\Lambda_b}) = 0$, while lower histograms correspond to $\sqrt{2}/2$.