Polarized beam studies at Budker Institute

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Abstract. This report describes activities in the polarization field at Budker Institute of Nuclear Physics. It concerns the radiative beam polarization, flipper and Siberian snake techniques and accurate beam energy measurements for electron and proton machines. Three new developments in the field of radiative polarizations are done recently at BINP. Filtering mechanism for polarized antiprotons obtaining is discussed. An general approach to consideration of RF-fields influence on spin dynamics at electron (positron) and hadron accelerators is given.

1. Introduction
Polarized beams is the traditional theme in Budker Institute of Nuclear Physics scientific programm starting from 60-th years of the last century. It’s good known the general approach to a description of spin motion in accelerators, developed in BINP at earlier years and, as a consequence of that, a set of theoretical and experimental studies radiative beam polarization.

Let’s remind the BMT equation, which describes spin motion in electromagnetic fields: [1]

\[
\frac{d\mathbf{S}}{dt} = \dot{\mathbf{S}} = [\mathbf{\Omega} \times \mathbf{S}] \\
-\mathbf{\Omega} = \left(\frac{q_0}{\gamma} + q'\right)\mathbf{B}_\perp + \frac{q_0 + q'}{\gamma} \mathbf{B}_\parallel + \left(\frac{q_0}{\gamma + 1} + q'\right)[\mathbf{E} \times \mathbf{V}],
\]

where \(q_0\) and \(q'\) are normal and anomalous parts of the particle gyro-magnetic ratio; \(\mathbf{B}_\perp\) and \(\mathbf{B}_\parallel\) are magnetic field components along and transverse to the particle velocity \(\mathbf{V}\). Since in circular accelerators the one-turn energy change is relatively small, we can neglect, in the first approximation, the electric field: \(\mathbf{E} = 0\). Following the usual approach for orbital motion, we use as the independent variable the generalized azimuth \(\theta\) and subdivide the spin precession vector in two parts: \(\mathbf{\Omega} = \mathbf{W}_0(\theta) + \mathbf{w}(\theta)\), where \(\mathbf{W}_0(\theta)\) contains only fields on the Closed Orbit \(R_0\), while \(\mathbf{w}(\theta)\) denotes all of the other terms (contributions from closed orbit imperfection and orbital oscillations). One can treat \(\mathbf{w}\) as a small perturbation for the spin motion.([2]–[4]) In the accelerator vector triad \(\mathbf{e}_x, \mathbf{e}_y = \mathbf{V}/V, \mathbf{e}_z = [\mathbf{e}_x \times \mathbf{e}_y]\) components of the precession vector \(\mathbf{W}_0\) can be presented in the next form: \(^1\)

\(^1\) We use dimensionless units: fields are normalized to mean guiding field \(B_0 = 1/2\pi \int B_z d\theta\); length and time are measured correspondent in units of mean radius \(R\) and revolution time.
\[ W_0x = \nu_0 K_x; \quad K_x = \frac{B_x}{B_0}; \]
\[ W_0y = (1 + a) K_y; \quad K_y = \frac{B_y}{B_0}; \]
\[ W_0z = \nu_0 K_z; \quad K_z = \frac{B_z}{B_0}, \]

(2)

where we introduce the particle magnetic anomaly \( a = q'/q_0 \) and denote \( \nu_0 = \gamma \cdot a \). On the reference orbit the equation has one solution \( n_0 \), which is periodic around the ring, i.e. \( n_0(\theta + 2\pi) = n_0(\theta) \). There are also two other linearly independent solutions: the orthogonal complex vectors \( \eta \) and \( \eta^* \), which rotate around \( n_0 \) with spin tune \( \nu \): \( \eta(\theta + 2\pi) = e^{i2\pi \nu} \eta(\theta); \quad \eta^*(\theta + 2\pi) = e^{-i2\pi \nu} \eta^*(\theta) \).

We can introduce also for each momentum-off particle an precession axis \( n \), which depends on a particle trajectory. This axis is periodical versus all phases of the orbital motion. It’s clear, that in the linear approximation \( n \) can be expressed as a combination of \( n_0 \) and \( \eta^*(\theta) \): \( n(\theta) = \sqrt{1 - C^2} n_0(\theta) + Re(iC \eta^*(\theta)). \)

Since \( w \ll W \), complex constant \( C \) is found by the perturbation theory:
\[ C' = \sqrt{1 - C^2} w_\perp - i w_\parallel, \]

(3)

where \( w_\perp = w \cdot \eta^* \) and \( w_\parallel = w \cdot n_0 \).

2. Radiative polarization

Such approach has leaded to a creation of the radiative polarization theory for real machines. According this theory particle spins build up along \( n_0 \) with a “damping” time \( \tau_p \), and finally achieve an equilibrium polarization degree \( P_{eq} \), which are described by so called DK formulae:[4]

\[ P_{eq} = -\frac{8}{5\sqrt{3}} \frac{\alpha_-}{\alpha_+}; \quad \tau_p^{-1} = \frac{5\sqrt{3} e\hbar \gamma^5}{8 m^2 c^3} \alpha_-; \]

\[ \alpha_- = \langle \rho^{-3} b \cdot (n_0 - d) \rangle; \quad \alpha_+ = \left\langle |\rho|^{-3} \left[ 1 - \frac{2}{9} (n_0 \cdot V)^2 + \frac{11}{18} d^2 \right] \right\rangle, \]

(4)

where \( b = \frac{B}{|B|} \) is unit vector along the guiding field, \( \rho \) is a curvature radius and \( d = \gamma \frac{\partial n}{\partial \gamma} \) is a periodical spin-orbit coupling vector, which shows a sensitivity of spin precession axis \( n \) to an energy change. Here and later the angle brackets \( \langle \rangle \) denote an average over time and a particle ensemble in the phase space.

First summand in \( \alpha_- \) (4) is responsible for a direct instant spin rotations by magnetic moment radiation fields. An effect of such self polarization has been predicted by Sokolov and Ternov in 1963 year for homogeneous magnetic field.[5] From that time this mechanism is good studied, because it takes place at flat machines, where \( n_0 \parallel b \).

Recently at BINP, the radiative polarization about \( \simeq 0.6 \) was measured at VEPP-2000 storage ring. This machine lattice has four strong solenoids with field up to 13 T, which is used as final focus system. Here the spin precession vector \( n_0 \) (vertical in the arcs) is rotated by each solenoid by 90 degrees.
The second term in $\alpha_\perp$ arises from an action on a particle trajectory by the same radiation fields of magnetic moment. By that, the vector $n$ is kicked randomly versus $n_0$. A final result depends on the projection of the spin orbit coupling vector $d$ on the guiding field direction $b$. That kinetic polarization mechanism takes place at storage rings with one Siberian snake, when $n_0$ rotates in the horizontal plane with spin tune $\nu = 1/2$ and $d \parallel b$.[6] There is one more advantage at machines equipped by one snake. Polarization directs along the snake axis at opposite to the snake insertion azimuth.[7] Two storage rings AmPS and SHR have worked in such scheme (see Fig.1). But polarization time at both rings was much longer, than beam life time at the beam energy about 1 GeV. A number of experiments with longitudinal polarized beam and polarized internal targets have been done. But polarized beam have been injected from polarized electron sources.

A simple analysis of the expression (4) on a maximum polarization gives $P_{eq} \simeq 0.8$, when $|d| \simeq 1.2$. But in reality the spin-orbit coupling vector strong changes around a machine. For example, a behaviour of $|d(\theta)|$ for SHR storage ring is demonstrated at the Fig.2.

![Figure 1. Layout of SHR.](image1)

![Figure 2. Absolute value of $d$ along the orbit.](image2)

The considerable grow $|d|$ at one half of the ring results to low polarization level $P_{eq} \simeq 0.1 - 0.2$. To improve the kinetic polarization Korostelev and Shatunov have proposed at parts of the ring, where $|d| \sim 1.2$, to install (see Fig.3) two asymmetric wigglers with 10 T field in the central part (10 cm) and low field parts at left and right sides.[8] As it’s shown at the Fig.4, wiggler insertions lead to enhancement of polarization up to $P_{eq} \simeq 0.6$ at the energy $E=1$ GeV. Unfortunately, a study of the kinetic polarization as well as a realization of this proposal have never done, because AmPS and SHR were closed around 2000 year.

However, a new interest to the kinetic polarization appears recently in connection with new projects of super C-tau and super B-factories. To arrange condition for polarization and save the longitudinal polarization in one Interaction Point, it’s enough to install odd number of solenoidal Siberian snakes. By that, one of them has to be located at opposite to the IP straight section. One of possible scheme is presented by I Koop on this Symposium.[9] Project luminosity of the Novosibirsk C-tau factory is about $1 \cdot 10^{35}$ cm$^{-2}$ sec$^{-1}$. There are two scenarios to obtain the polarization and don’t loose in the luminosity. At the moment, technics of the polarized electron sources is very high developed. On the other hand, there are a number of processes, when only one polarized beam is enough to extract spin depended physics. For such experiments polarized electrons can be delivered by a linear accelerator directly from the source with polarization about $P \simeq 0.9$. Beam life time due to high luminosity is much shorter, than the polarization time without special wigglers. Finally, one can get with frequent
injection high polarized electron beam and unpolarized positrons. The Fig. 5 shows results of calculations for one, three and five Siberian snakes at the electron ring.

But, there are also ideas of experiments with two polarized beams. In this case, we have to equip the positron storage ring (or both rings) by snakes and asymmetrical wigglers, while the polarization time will be shorter, than beam life time.

![Figure 3. Asymmetric wiggler for SHR.](image1)

**Figure 4.** Equilibrium polarization degree with wiggler “on”.

![Figure 5. Equilibrium longitudinal polarization at C-τ factory versus beam energy.](image2)

**Figure 5.** Equilibrium longitudinal polarization at C-τ factory versus beam energy.

3. Spin control by RF-fields
Pioneering works with application of RF-field for spin manipulations has opened way to precise spin tune measurement as well as spin flipping technics.[10] A number of precise experiments have been done at BINP’s colliders VEPP-2M and VEPP-4 using RF-field devices — depolarizer
and flipper. However, till now it occurs some misunderstandings of the RF-field influence on spin motion. So, we give shortly here an introduction in the problem.

We restrict consideration to planar machines and start with a proton ring, where the spin tune is $\nu = \nu_0$ and $n_0 = e_2$; $\eta = (e_x - ie_y) e^{i \theta}$. Let’s apply on a short piece of the orbit $\Delta l$ a longitudinal RF-field $\tilde{K}_y \cos(\nu_0 \theta)$ (RF-solenoid) with a frequency $\nu_0$, which is an external spin perturbation. It can be presented as a number of circular harmonics $w = \sum_k w_k e^{i(k + \nu_0) \theta}$ with equal amplitudes $w_k = (1 + a) \tilde{K}_y \frac{\Delta l}{2\pi}$.

A more complicated situation occurs for the application of a radial RF-field $\tilde{K}_x \cos(\nu_0 \theta)$ (RF-dipole). This field disturbs not only spin motion, but also excites also forced vertical oscillations. As a result, spin gets additional kicks from off-orbit fields in dipole and quadrupole magnets around the machine. The tight frame of this paper doesn’t allow a mathematical description of this process and we refer readers to papers devoted to this topic.\[11]-[12] The linear spin response formalism has been developed for simultaneous treating of the orbit and spin dynamics. Based on this formalism the code “ASPIRIN” ([14]) calculates 5 complex response functions $F_1(\theta) - F_5(\theta)$ which satisfy the periodicity condition: $|F_i(\theta + 2\pi)| = |F_i(\theta)|$. These response functions characterize the sensitivity of the spin precession axis $n$ to kicks of orbital variables $(X^T = (p_x, x, p_z, z, \delta \gamma/\gamma, 0))$ for specific machines optics and working points. In the case of an ideal flat machine, $\frac{\partial p_c}{\partial \nu_c} = \nu_0 \Re (i F_3(\theta) \eta^*)$. Thereby the RF-dipole, applied at azimuth $\Theta$, creates a set of perturbing harmonics with amplitudes $w_k = \nu_0 \tilde{K}_y F_3(\Theta) \frac{\Delta l}{2\pi}$.

Next, it’s necessary to take into account synchrotron oscillation of the particle energy, inherent to accelerators: $\gamma = \gamma_0 + \Delta \gamma \cos(\nu_0 \theta)$. Usually, the synchrotron tune $\nu_0 \sim 10^{-2} - 10^{-3}$ is much less than other orbital frequencies. Hence, the spin tune is modulated by the synchrotron tune. It results in a modification of the RF-field spectrum. Each line of the spin perturbation is transformed to a set of side bands resonances $k \pm \nu_0 \pm m \nu_s$. The side band harmonics are given by the expression:\[2]

$$|w^m_k| = |w_k| J_m(\lambda),$$

where $J_m(\lambda)$ are the Bessel functions. As a rule, the modulation index $\lambda = \frac{\nu_0}{\nu_s} (\Delta \gamma)$ $\ll 1$, at proton and electron accelerators. For RF spin resonances not all sidebands overlap ($w^m_k \ll \nu_s$). It’s important to emphasize that the central line of the spectrum corresponds to the mean energy $\gamma_0$ of the beam. Moreover, the width of this spectrum line, averaged over the synchrotron oscillations, shrinks to a size of $\sigma_\nu \sim (\Delta \gamma/\gamma)^2$.\[18] Assuming a Gaussian particle distribution in the longitudinal phase space, mean value of the central resonance strength ($w_0^0 \equiv w$) and it’s rms deviation $\sigma_w$ are given by the expressions:

$$w^2 = w^2_0 I_0(\Lambda) e^{-\Lambda}; \quad \sigma_w^2 = \frac{w^2_0}{2} I_1(\Lambda) e^{-\Lambda},$$

where $I_0, I_1$ are the modified Bessel functions from the argument $\Lambda = \frac{\nu_0}{\nu_s} \sigma_\gamma; \left(\frac{\sigma_\gamma^2}{\gamma} = \frac{\Delta \gamma^2}{\gamma^2}\right)$.

Evidently, spin control operations have to be done at the central resonance. The choice of frequency $\nu_0$ is determined by concrete conditions of an experiment, but always one must satisfy the resonance condition: $|\nu_0 \pm k - \nu_0| \ll 1$. It’s more visual to consider this situation in a frame rotating at the resonant frequency, where the spin motion is simply precession in a “field” $h = e e_x + w e_y$, where $e = \nu_0 - \nu_0 \pm \frac{k}{2}$ is the resonance tune, see Fig.6. The precession frequency $h = \sqrt{e^2 + w^2}$ in the resonance case ($e = 0$) decays to $w_k$, which can be taken as a strength of the resonance.

The spin resonance crossing was described by Froissart-Stora [16], who found that in result of passing from $\epsilon = \infty$ up to $\epsilon = -\infty$ with a tune rate $\epsilon = \text{const}$ the residual projection averaged
over spin phases $S_z = \langle S \cdot n_0 \rangle$ is described by a formula:

$$S_z = S_z(0) \left( 2 e^{-\frac{\pi w^2}{2 \dot{\epsilon}}} - 1 \right). \quad (7)$$

The final result depends on the spin phase advance near the resonance center: $\Psi = \int h dt \sim w^2/\dot{\epsilon}$. The polarization is preserved by an adiabatic change of parameters ($\Psi \gg 1$) and flips down together with the $n = h/h$-axis. In the opposite case of a fast crossing ($\Psi \ll 1$), the spin only tilts slightly from its initial direction with a depolarization $\Delta S_n \approx \frac{w^2}{\dot{\epsilon}}$. Polarization losses may attain \sim 100% in intermediate situations $\Delta \phi \leq 1$.

At electron accelerators there are other limitations for RF-device usage. Quantum fluctuations of the synchrotron radiation lead to so called, “spin diffusion”, which is enhanced at spin resonances.[3] In a general case, a quantum emission in the resonant frame brings at the same instant jumps of the resonance tune ($\delta \epsilon$), the resonant harmonic ($\delta w_k$) and it’s phase ($\delta \phi$), see Fig.7. At that, the precession axis $n = h/h$ gets a kick $\delta n$; But spin vector does not change (we neglect here “spin-flip” quantum emissions). However, the projection $S_n$ undergoes a change $\delta S_n = -1/2 (\delta n)^2 S_n$. A consequence of random kicks results in a polarization loss with a decay time $\tau_d = 1/2 \langle \frac{d}{dt} (\delta n)^2 \rangle$.

$$\langle (\delta n)^2 \rangle = \left( \delta \arctan \left( \frac{|w_k|}{\epsilon} \right) \right)^2 + \frac{w_k^2}{\epsilon^2} (\delta \phi_k)^2, \quad (8)$$

In the case of the RF-resonance ($\delta \phi = 0$), we assume the time of the resonance crossing $\Delta t \sim w/\dot{\epsilon}$ is much longer, than the radiative damping time $\tau_0$ and the characteristic times of the orbital motion. Therefore we can consider, in average, numerous jumps of $\delta \epsilon$ and $\delta w$ only as a diffusion around mean values $\bar{\epsilon}$ and $\bar{w}$ with corresponding rms deviations $\sigma_{\epsilon}$, and $\sigma_w$. So, from (8) we find the depolarization time, caused by the quantum fluctuations:

$$\tau_d^{-1} = \frac{1}{2} \left( \frac{\epsilon \delta w - w \delta \epsilon}{w^2 + \epsilon^2} \right)^2 \simeq \frac{\tau^2 \sigma_{\epsilon}^2 + \bar{w}^2 \sigma_w^2}{(\bar{w}^2 + \tau^2)^2} \cdot \tau_0^{-1}. \quad (9)$$
To obtain above expression, we ignored the interference term by the averaging over time \((\delta \epsilon \cdot \delta w = 0)\) and used usual formulas for an equilibrium state: 
\[
\sigma^2_{\epsilon} = 1/2 \left\langle \frac{d}{dt} (\delta \epsilon)^2 \right\rangle \cdot \tau_0; \\
\sigma^2_{w} = 1/2 \left\langle \frac{d}{dt} (\delta w)^2 \right\rangle \cdot \tau_0 \\
\text{with } \sigma_{\nu} \text{ from (6), where } \sigma^2_{\gamma} = 1/2 \left\langle \frac{d}{dt} (\delta \gamma)^2 \right\rangle \cdot \tau_0.
\]

At the next step we employ the FS-formula for electron machines, taking into account the spin diffusion under the condition of adiabatic crossing \((\psi \gg \frac{\nu}{\sqrt{2}}\)). We calculate the beam polarization \(\zeta(t) = \langle S_n \rangle = \zeta(0) \cdot \int_0^t e^{-t/\tau} dt\), while resonance crossing with different \(w\) and \(\dot{\epsilon}\).

For example, we use parameters of VEPP-2M storage ring \((E=700 \text{ MeV})\). The initial polarization is always \(\zeta(0) = 1\). Fig.8 shows three curves \(\zeta(t)\) versus \(\epsilon(t)\) for different RF resonance amplitudes and spin tune spreads but for the same crossing rate \(\dot{\epsilon} \cdot f_0 = 400 \text{ Hz/sec}\): (solid line) \(w = 1 \cdot 10^{-5}; \sigma_{\nu} = 3 \cdot 10^{-7}\); (dotted line) \(w = 4 \cdot 10^{-5}; \sigma_{\nu} = 1 \cdot 10^{-6}\); and (dashed line) \(w = 1 \cdot 10^{-4}; \sigma_{\nu} = 3 \cdot 10^{-6}\). These calculations demonstrate clearly the influence of the spin diffusion and spin tune spread. Even increasing the RF amplitude by ten times does not help to avoid some depolarization, when the spin tune spread grows up to three times. So, using such “simulations”, it’s possible to choose RF device parameters for successful spin flip. Moreover, the measurement of a residual polarization in the case of a reasonable polarization loss appears as a way to measure the spin tune spread and minimize it, if that is necessary.[18]

![Figure 8. Spin flip by RF-field.](image8.png)

![Figure 9. Beam resonant depolarization.](image9.png)

It’s interesting also to study the opposite case of a small RF amplitude. Fig.9 presents three other curves, where we fixed the spin tune spread \(\sigma_{\nu} = 1 \cdot 10^{-6}\) and the crossing rate \(\dot{\epsilon} \cdot f_0 = 2 \text{ Hz/sec}\), but changed the amplitude \(w\): (solid line) \(w = 2 \cdot 10^{-7}\); (dotted line) \(w = 5 \cdot 10^{-7}\) and (dashed line) \(w = 1 \cdot 10^{-6}\). One can see from Fig.9 the resonant depolarization by the RF-field. Decreasing the RF power provides a measurement of the spin tune with accuracy up to its spread \(\sigma_{\nu}\). In turn, the spin tune determination is simultaneously the absolute mean energy measurement, since the magnetic anomalies are well known. Beam energy calibration has been routinely used at electron-positron colliders in precise experiments for secondary particle mass measurements.[19] Recently, the beam energy was measured at VEPP-4M with accuracy \(\approx 10^{-7}\). This result is restricted already by a knowledge of electron mass, which is known only in three times better. The coherent spin rotation by 90 degrees and full spin flip were crucially important at VEPP-2M in the experiment for electron and positron anomalous magnetic moments comparison.[20]

Till now, RF spin flip has been studied for protons and deuterons at the IUCF storage ring (see, for instance [21]) and at the synchrotron COSY.[22] At first, the authors pointed out spin flip by RF-fields was going, as a rule, with very low polarization losses \((\approx 10^{-2} \text{ per pass})\). At both machines the experiments were carried out with RF-solenoids and RF-dipoles.
and the authors claimed to find “unexpected discrepancies” between the measured values and “the theoretical expressions for the spin flip resonance widths”. However, recent analysis of the COSY data, based on the formalism of the spin response functions, has found an amazing accordance of the experiment and the theory as for both kinds of the flippers, for both types of particles.[12] The main reason for the “unexpected discrepancies” is explained by Fig. 10, which presents results of ASPIRRIN calculations of $|F_3|$ values along the COSY orbit for protons and deuterons, that are distinguished in more than 1000 times in the point of the RF flipper.

![Graph](image1.png)

**Figure 10.** $|F_3|$ along the COSY orbit for protons and deuterons.

It’s necessary to pay attention to RF-field control at machines with Siberian snakes. Many of main principles for that were worked out by Koop and Shatunov.[23] Since spin tune at this case $\nu = \frac{1}{2}$, hence the location of the spin resonance $\nu_d = \nu$ and resonance $1 - \nu_d = \nu$ coincide, so the condition for an isolated spin resonance does not satisfied. Hence the Snake should be detuned, to shift the spin tune slightly from $\nu = \frac{1}{2}$, so the spin flip resonance and it’s mirror are separated. For the solenoidal RF-field the resonance strength does not change: $w_k = (1 + a)K_0 \frac{\Delta l}{4\pi R}$. But for the transverse RF-fields (AC dipoles) situation differs due to the forced by AC dipole betatron oscillations. Presence of the Siberian snakes (and other spin rotators) modify the spin response functions $F_1$ or $F_3$. Mane and Shatunov carried out an analysis for different combinations of the Siberian snakes and spin rotators at RHIC.[24] They found, that, in spite of the $F_3$ modification, the contribution of the forced betatron oscillation to the spin flip resonance strength is not negligible in contrast to assumption in Mai and Roser’s paper.[25] However, an idea, formulated in [25], to compensate the mirror harmonics by an additional AC-dipole is promising. Taking into account the spin response function formalism, a scheme with two AC dipoles is able to make the spin flip with perfect full Siberian snake ($\nu = \frac{1}{2}$) and arbitrary spin rotators. By that, an required power supply for AC dipoles will be very moderate.

4. Related activities
On this Symposium tree more talks are given from BINP community. I Rachek and V Stibunov present results of experiments with polarized hydrogen and deuterium internal target at the storage ring VEPP-3. Atomic beam source at VEPP-3 provides a flux $\cong 8 \times 10^{16}$ atoms per second. They have studied a number of processes by electrons scattering on the nucleons: measurements of elastic deuteron form-factors ($G_C$ and $G_Q$) and quasi-elastic proton knock-out. For the first time a pions photo-production ($\gamma + d \rightarrow d + \pi^0$) and deuteron photo-disintegration ($\gamma + d \rightarrow p + n$) have been observed.
Other interesting topics has been investigated by group of BINP’ theorists in a connection with plans to construct a polarized proton-antiproton collider in the frame of FAIR project. Few year ago A Milstein and V Strakhovenko have found a right explanation of the filtering experiment for protons scattered on polarized internal target.[26] Later, antiprotons polarization by the filtering mechanism have estimated in the Paris, Juelich and Nijmigen potential models. S Salnikov summarizes these results at his talk on the Symposium. In spite of different results obtained in each model for $P_\perp$ and $P_\parallel$, the polarization effect is quite promising. But, of course a test experiment is needed.

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