Symmetry of crystals with diamond structure and formation of radiation defects induced by negative muon

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Abstract. Positive and negative muons ($\mu^+$ and $\mu^-$) are widely used for study condensed matter at present. Negatively charged muon is captured by nuclei rapidly at 1s-state and forms $\mu$-atom. Because of the muon mass is at 207 time greater than the electron mass, negatively charged muon screens effectively a nucleus with a charge $Z$ and so models a single atomic isotope with the charge equal to $Z - 1$. To interpret experimental results one supposes as usual that muons are thermalized in a condensed matter for a time shorter than an equipment resolution time. Experimental technique development showed that this suggestion is not true for some targets.

A neutralization process of a radiation defect induced by negative muon in crystals with diamond structure (Diamond, Silicon and Germanium) is considered in this work. Estimates of the velocity of neutral acceptor center ($\mu_A^0$) formation connected with radiation transitions are obtained. Radiation transitions take place when a neutral atom ($\mu_A^0$) changes its configuration and forms chemical links with nearest neighbours in hybrid states. Symmetry of crystals with diamond structure allows us to express the velocity of neutral acceptor center formation in terms of one matrix element only.

diamond structure, radiation defect, negative muon, radiation transition, hybridization, muonic atom

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1. Introduction

Nowadays positive and negative muons ($\mu^+$ and $\mu^-$) are widely used for different study of condensed matter, behaviour of hydrogen-like light impurities and modeling of chemical processes with atomic Hydrogen (see, e.g. [1]). $\mu^+$ and $\mu^-$ behaviour in a medium is strongly different: positively charged muon is a light impurity and models from a chemical point of view a “light hydrogen isotope”. Negatively charged muon is captured rapidly by a nucleus at 1s-state and forms the so called $\mu$-atom.

Because of the muon mass is at 207 time greater than electron mass, negatively charged muon screening effectively a nucleus with a charge $Z$. In this case the effective charge of the nucleus appears for 1 unit less than primary charge. So, this kind of radiation defect models a single atomic isotope with a charge equal to $Z - 1$. This result is well known for the initial stage of muonic research (see, e.g. [2]) and gives rise to the foundation of the muonic ($\mu$SR) method for matter study. Systematic research of impurity formation with nucleus charge $Z - 1$ was carried out in condensed matter at the early stages of $\mu$SR-researches [3, 4].

Study of acceptor centers in semiconductors with $\mu^-$ was suggested in the work [5]. For example, in Silicon ($Z = 14$) negative muon is captured by the nuclear and forms an “isotope” of Aluminium $\mu$Al with nuclear charge equals to $Z = 13$. This kind of impurity is in interest because it is one of the main acceptor impurities in silicon semiconductors. It was shown [6] - [8] that with negative muons we can get very valuable information on a hyperfine structure and interactions of acceptor centers with the lattice in different semiconductors.

To interpret experimental results one supposes as usual that muons are thermalized in a condensed matter for a time shorter than an equipment resolution time. So, this suggestion determines a model for a description of the muon spin polarization. Experimental technique development and broadening of types of objects in research show that this suggestion is not true for any targets. In particular, theoretical explanation of experimental results with positive muons in solid Hydrogen and Helium (see [9, 10]) and [11]) was successful after taking into account thermalization processes [12–14].

Recent experiments with negative muons in Diamond [15, 16] show that interpretation of the results failed in the suggestion of rapid ($< 10^{-10}c$) thermalization of forming the muonic acceptor center $\mu$B. Thus, examination of a a thermalization process of a muonic acceptor center as a radiation defect in crystals, is actual.

A neutralization process of a radiation defect induced by negative muon in crystals with diamond structure (Diamond, Silicon and Germanium) is considered in this work. Estimations of the velocity of neutral acceptor center ($\mu$A)$^0$ formation connected with radiation transitions are obtained (here A is the appropriate acceptor atom in the lattice).
2. States of neutral ($\mu A$)\textsuperscript{0} center in the lattice

Estimations show that a substitution impurity center formed after negative muon capture at a $K$-shell is neutralized rapidly capturing an appropriate number of electrons. A typical neutralization time is of the order of $t \sim 10^{-12}$ s and the number of electrons for example is 5 to Diamond, 13 to Silicon. Muonic impurity atom is in an excited state just after formation because it’s chemical bonds with host atoms are broken. In accordance with standard idea of quantum chemistry only electrons with the same general quantum number can create chemical bond if they were on an unfilled energy level of the atom. In this case they form hybridized states. For lattices with diamond structure electron states $ns$ and $np$ are represented with the equal probability, where $n = 2, 3$ and $4$ relate to C, Si and Ge respectively. Really, hybridized states are formed in atomic time, but chemical bonds formation is determined by exchange interactions that are smaller than coulomb interactions which forms appropriate atomic configuration.

When a chemical bond is formed a significant value of energy (of the order of some eV) can be emitted. In gasses and liquids this excess energy can be transferred to the third body. This kind of energy transfer in crystal must be connected with a phonon emission. One-phonon emission with the energy of $\geq 1$ eV in covalent crystal is impossible. So, a transfer of this energy value could be realized in the case of a multy-phonon process. This kind of processes have very small probability. Thus, a radiation transition with a photon emission seems to us more preferable with respect of the other processes.

Let’s suppose that an impurity atom with the nuclear charge $Z - 1$ is formed as a result of a neutralization process in an atomic time and has an atomic configuration where electrons at the external shell are in hybridized but not in the ground state:

$$ |\psi_{\text{in}}\rangle = |ns\; np^2\rangle. \quad (1) $$

We consider that the electron configuration with the main quantum number less then $n$ is completely filled and the state of such electrons is described by nonperturbed wave function of the free atom. The initial state of the radiation defect in a diamond lattice is sketched in Fig. 1a.

The hybridized state (1) forms a chemical bond, if it possesses by maximal spin $S = 3/2$, and a spacial part of its wave function one may represent in the view of equiprobable superposition of three Slater’s determinants:

$$ \Psi_{sp^2}(r_1, r_2, r_3) = \frac{1}{\sqrt{3}} (\Psi_{+1}(r_1, r_2, r_3) + \Psi_{0}(r_1, r_2, r_3) + \Psi_{-1}(r_1, r_2, r_3)), \quad (2) $$

where

$$ \Psi_M(r_1, r_2, r_3) = \frac{1}{\sqrt{6}} \begin{pmatrix} \psi_{ns}(r_1) & \psi_{np,m}(r_1) & \psi_{np,m'}(r_1) \\ \psi_{ns}(r_2) & \psi_{np,m}(r_2) & \psi_{np,m'}(r_2) \\ \psi_{ns}(r_3) & \psi_{np,m}(r_3) & \psi_{np,m'}(r_3) \end{pmatrix}. \quad (3) $$

Here are $M = m + m'$, $m, m' = 0, \pm 1$. 
Radiation defect $\mu B$ state in a diamond lattice: 

- all bonds with host atoms are broken in the initial state and the defect state is determined by the hybridized function $2s^22p^2$; 
- three electrons of the impurity form chemical bonds with host atoms in the final state and an unsaturated (broken) bond is equiprobable for four nearest neighbours of the cluster $(BC_4)^0$.

A spatial part of the wave function of the defect in a final state might be represented in the view of a superposition of three hybridized states forming the chemical bond with host atoms of the lattice:

$$\Psi_{Cl}(\mathbf{r}_1, \mathbf{r}_2, \mathbf{r}_3) = \frac{1}{\sqrt{6}} \sum_P (-1)^P \Psi_{\mu A}(\mathbf{P}(\mathbf{r}_1, \mathbf{r}_2, \mathbf{r}_3)).$$  \hspace{1cm} (4)

Here summation is carried out on all transpositions of the valence electrons of the impurity and

$$\Psi_{\mu A}(\mathbf{r}_1, \mathbf{r}_2, \mathbf{r}_3) = \frac{1}{2} \left( \psi_{n_1}(\mathbf{r}_1)\psi_{n_2}(\mathbf{r}_2)\psi_{n_3}(\mathbf{r}_3) + \psi_{n_2}(\mathbf{r}_1)\psi_{n_3}(\mathbf{r}_2)\psi_{n_1}(\mathbf{r}_3) + \psi_{n_3}(\mathbf{r}_1)\psi_{n_1}(\mathbf{r}_2)\psi_{n_2}(\mathbf{r}_3) + \psi_{n_1}(\mathbf{r}_1)\psi_{n_2}(\mathbf{r}_3)\psi_{n_3}(\mathbf{r}_2) + \psi_{n_2}(\mathbf{r}_1)\psi_{n_3}(\mathbf{r}_3)\psi_{n_1}(\mathbf{r}_2) + \psi_{n_3}(\mathbf{r}_1)\psi_{n_1}(\mathbf{r}_3)\psi_{n_2}(\mathbf{r}_1) \right).$$  \hspace{1cm} (5)

The unit vectors $\mathbf{n}_a$ are directed from the impurity to nearest neighbours (along the directions of the chemical bonds). The one-particle functions $\psi_{n_a}(\mathbf{r})$ are the hybridized states with directed bonds:

$$\psi_{n_a}(\mathbf{r}) = \alpha_a \psi_{n_a}(\mathbf{r}) + \beta_{n_a} \psi_{n_p,n_a}(\mathbf{r})$$  \hspace{1cm} (6)

and satisfy normalized conditions:

$$\int \psi_{n_a}^*(\mathbf{r})\psi_{n_b}(\mathbf{r})d\mathbf{r} = \delta_{ab}. \hspace{1cm} (7)$$

From this condition we can see relations for the coefficients in the superposition (6):

$$\alpha_a^2 + \beta_{n_a}^2 = 1, \quad \sum_a \alpha_a^2 = 1, \quad \sum_a \beta_{n_a}^2 = 3, \quad \alpha_a \alpha_b + \beta_{n_a} \beta_{n_b} \cos(\mathbf{n}_a,\mathbf{n}_b) = 0. \hspace{1cm} (8)$$
3. Formation of neutral center \((\mu, A A_4)^0\)

A life-time of the exited state (1) is determined by a rate of a radiation transition in a bond state and might be calculated with Fermi’s “golden rule” relation:

\[
\frac{dw}{dt} = \frac{2\pi}{\hbar} |\langle \psi_{in} | \hat{V}_{\text{rad}} | \psi_{cr} \rangle|^2 \delta(E_f - E_i)d\nu_f.
\]  

(9)

The interaction operator one may write in a form of:

\[
\hat{V}_{\text{rad}} = \frac{e}{m\epsilon} \sum_a \hat{p}_a A(r_a),
\]  

(10)

where \(A(r)\) is the vector-potential of the free radiation field.

Let’s consider now only the term for one electron with \(a = 1\) in the operator (10) to simplify the following calculations. In this case matrix elements of the perturbation operator could be represented by the expression in the form of

\[
\langle \psi_{in} | \hat{p}_1 | \psi_{cr}\rangle_{sa} = \int \psi_{ns}^*(r_1) \hat{p}_1 \psi_{s}^*(r_1) dr_1 \int \psi_{np,m}^*(r_2) \psi_{s}^*(r_2) dr_2 \int \psi_{np,m'}^*(r_3) \psi_{s}^*(r_3) dr_3
\]  

(11)

and

\[
\langle \psi_{in} | \hat{p}_1 | \psi_{cr}\rangle_{pa} = \int \psi_{np,m}^*(r_1) \hat{p}_1 \psi_{s}^*(r_1) dr_1 \int \psi_{ns}^*(r_2) \psi_{s}^*(r_2) dr_2 \int \psi_{np,m'}^*(r_3) \psi_{s}^*(r_3) dr_3.
\]  

(12)

where indexes are \(a, b, c = 1, 2, 3, 4\).

If we direct the axis \(z\parallel n_1\), then the other three \(p\)-states in hybridized states (6) turn out as a result of rotation of the state \(\psi_{n1,0}(r)\) in the state with a rotation moment projection equals to zero on the axes \(n_b\). In this case we get the opportunity to calculate easy integrals incoming in the expressions (11) and (12):

\[
\int \psi_{ns}^*(r) \psi_{s}^*(r) dr = \alpha_a \langle ns|ns \rangle = \alpha_a;
\]  

(13)

\[
\int \psi_{np,m}^*(r) \psi_{s}^*(r) dr = \beta_a \langle np, m|\hat{R}(\theta_a, \varphi_a)|n1, 0 \rangle,
\]  

(14)

where \(\hat{R}(\theta_a, \varphi_a)\) is the rotation operator and the matrix elements (14) are determined by the second column of the rotation matrix:

\[
\langle n1, \pm 1|\hat{R}(\theta_a, \varphi_a)|n1, 0 \rangle = \mp \frac{1}{\sqrt{2}} \sin \theta_a e^{-i\varphi_a}, \quad \langle n1, 0|\hat{R}(\theta_a, \varphi_a)|n1, 0 \rangle = \cos \theta_a.
\]  

(15)

We don’t reduce the similarity if consider a matrix element for the \(z\)-projection of the momentum operator. So, we have:

\[
\int \psi_{ns}^*(r_1) \hat{p}_{1z} \psi_{s}^*(r_1) dr_1 = \beta_a \int \psi_{ns}^*(r_1) \hat{p}_{1z} R(\theta_a, \varphi_a) \psi_{n1,0}(r_1) dr_1 = \beta_a \cos \theta_a I_{sp}^{(n)},
\]  

(16)

\[
\int \psi_{np,m}^*(r_1) \hat{p}_{1z} \psi_{s}^*(r_1) dr_1 = \alpha_a \int \psi_{n1,0}^*(r_1) \hat{p}_{1z} \psi_{ns}(r_1) dr_1 = \alpha_a I_{sp}^{(n)},
\]  

(17)

where

\[
I_{sp}^{(n)} = \int \psi_{ns}^*(r) \hat{p}_{1z} \psi_{n1,0}(r) dr.
\]  

(18)

Let’s represent some intermediate calculations to clarify the understanding. The total matrix element in the expression (9) consists of 72 different items corresponding
to different matrix elements between states of the superpositions (2) and (4). However, it is enough to calculate only four of them. We are giving them below.

\[
\langle \Psi_{+1} | \hat{p}_{1z} | \Psi_{n_1,n_2,n_3} \rangle = -\frac{1}{4\sqrt{3}} \beta_3^2 \beta_4^2 \sin \theta_2 \cos \theta_3 e^{-i\varphi_2} + \cos \theta_2 \sin \theta_3 e^{-i\varphi_3} I_{sp}^{(n)} + \alpha \beta_3 \sin \theta_2 e^{-i\varphi_2} - \alpha \beta_3 \sin \theta_3 e^{-i\varphi_3} I_{sp}^{(n)\star}.
\]

The rest three matrix elements are determined by the other possible sets of \(n_a\) for the electrons with coordinates \(r_1, r_2\) and \(r_3\):\[
\langle \Psi_{+1} | \hat{p}_{1z} | \Psi_{n_2,n_3,n_4} \rangle = -\frac{1}{4\sqrt{3}} \beta_3^2 \beta_4^2 \sin \theta_2 \cos \theta_3 e^{-i\varphi_4} + \cos \theta_2 \sin \theta_3 e^{-i\varphi_4} I_{sp}^{(n)} - \alpha \beta_3 \sin \theta_2 e^{-i\varphi_4} - \beta_3 \alpha \sin \theta_3 e^{-i\varphi_3} I_{sp}^{(n)\star};
\]
\[
\langle \Psi_{+1} | \hat{p}_{1z} | \Psi_{n_3,n_4,n_1} \rangle = -\frac{1}{4\sqrt{3}} \beta_3^2 \beta_4^2 \sin \theta_2 \cos \theta_3 I_{sp}^{(n)} - \alpha \beta_3 \sin \theta_2 e^{-i\varphi_4} - \beta_3 \alpha \sin \theta_3 e^{-i\varphi_3} I_{sp}^{(n)\star} \sin \theta_4 e^{-i\varphi_4}.
\]
\[
\langle \Psi_{+1} | \hat{p}_{1z} | \Psi_{n_4,n_1,n_2} \rangle = \frac{1}{4\sqrt{3}} \beta_3^2 \beta_4^2 \sin \theta_2 \cos \theta_3 I_{sp}^{(n)} - \alpha \beta_3 \sin \theta_2 e^{-i\varphi_4} - \beta_3 \alpha \sin \theta_3 e^{-i\varphi_3} I_{sp}^{(n)\star} \sin \theta_4 e^{-i\varphi_4}.
\]

It is easy to see that transpositions of electrons in the superposition (5) don’t change expressions for the matrix elements (19)–(22). So, the number of transpositions in the state (5) with similar expressions reduces the total number of items for 6 times.

We get very cumbersome expression for the arbitrary values of the parameters \(\alpha, \beta, \theta\) and \(\varphi\). However, it is necessary to take into account that the system under consideration has a symmetry at less \(C_{3v}\). In this case the result could be essentially simplified. Let’s examine the simplest case at first, when the system has a tetrahedral symmetry and all parameters in the hybridized states (6) are equal to each other:

\[
\alpha = \alpha = \frac{1}{2}, \quad \beta = \beta = \frac{\sqrt{3}}{2}.
\]

If the vector \(n_4\) lies in the \(xz\) plane the angles \(\theta, \varphi\) are equal:

\[
\theta_1 = \varphi_1 = 0, \quad \theta_2 = \varphi_2 = \theta_4 = \theta, \quad \cos \theta = -\frac{1}{3}, \varphi_2 = -\varphi_3 = \frac{2\pi}{3}, \varphi_4 = 0.
\]

In the case of symmetrical structure we get the following expressions for the matrix elements:

\[
\langle \Psi_{+1} | \hat{p}_{1z} | \Psi_{n_1,n_2,n_3} \rangle = \frac{1}{8} \beta \left( \beta^2 \sin 2\theta I_{sp}^{(n)} - \alpha^2 \sin \theta I_{sp}^{(n)\star} \right),
\]

\[
\langle \Psi_{+1} | \hat{p}_{1z} | \Psi_{n_2,n_3,n_4} \rangle = \frac{3}{8} \left( 1 - \frac{1}{\sqrt{3}} \right) \beta \sin \theta \left( \beta^2 \cos^2 \theta I_{sp}^{(n)} - \alpha^2 I_{sp}^{(n)\star} \right),
\]

\[
\langle \Psi_{+1} | \hat{p}_{1z} | \Psi_{n_3,n_4,n_1} \rangle = -\frac{1}{4\sqrt{3}} \beta \sin \theta \left( \beta^2 \cos \theta I_{sp}^{(n)} - \alpha^2 I_{sp}^{(n)\star} \right),
\]

\[
\langle \Psi_{+1} | \hat{p}_{1z} | \Psi_{n_4,n_1,n_2} \rangle = \langle \Psi_{+1} | \hat{p}_{1z} | \Psi_{n_3,n_4,n_1} \rangle.
\]

Adding up the expressions (25)–(28), we get

\[
\langle \Psi_{+1} | \hat{p}_{1z} | \Psi_{C_4} \rangle = -\frac{1 - i\sqrt{3}}{8\sqrt{3}} \beta^3 \sin 2\theta \frac{\sin^2 \theta}{2} I_{sp}^{(n)}.
\]
The matrix elements for the state with $M = -1$ are calculated by the similar way:

$$\langle \Psi_{-1} | \hat{p}_{1z} | \Psi_{Cr} \rangle = - \frac{1 + i \sqrt{3}}{8 \sqrt{3}} \beta^3 \sin 2\theta \sin^2 \frac{\theta}{2} I_{sp}^{(n)}.$$ (30)

For the state with $M = 0$:

$$\langle \Psi_{0} | \hat{p}_{1z} | \Psi_{Cr} \rangle = - \frac{i}{2 \sqrt{2}} \beta^3 \sin^2 \theta \cos^2 \frac{\theta}{2} I_{sp}^{(n)}.$$ (31)

After substitution the values of the parameters (23)–(24) and adding up the expressions (29)–(31) we have:

$$\langle \Psi_{in} | \hat{p}_{1z} | \Psi_{Cr} \rangle = - \frac{\beta^3}{4 \sqrt{3}} \left( \sin 2\theta \sin^2 \frac{\theta}{2} + i \sqrt{2} \sin^2 \theta \cos^2 \frac{\theta}{2} \right) I_{sp}^{(n)} = \frac{1 - i}{72 \sqrt{2}} I_{sp}^{(n)}.$$ (32)

For the calculation of the integral $I_{sp}$ we take the appropriate wave functions of the hydrogen-like atom with an effective nuclear charge equal to $\tilde{Z}$.

In a diamond crystal a Boron $\mu$-atom $\mu_B$ is formed. It has the main quantum number $n = 2$ and we get for $\mu_B$

$$I_{sp}^{(2)} = -i \tilde{Z}, \quad \text{or in dimensional units } I_{sp}^{(2)} = -i \frac{\tilde{Z} mc^2}{h} = -i \frac{\tilde{Z} h}{a_0},$$ (33)

where $a_0$ is the Bohr radius.

In a silicon crystal an Aluminium $\mu$-atom $\mu_Al$ is formed. It has the main quantum number $n = 3$ and appropriate calculations for $\mu_Al$ give the following results:

$$I_{sp}^{(3)} = -\frac{10}{3 \sqrt{6}} \tilde{Z}, \quad \text{in dimensional units } I_{sp}^{(3)} = -i \frac{10}{3 \sqrt{6}} \frac{\tilde{Z} h}{a_0},$$ (34)

In a Germanium lattice a Gallium $\mu$-atom $\mu_Ga$ with the main quantum number $n = 4$ must be formed. The unknown value of the matrix element for $\mu_Ga$ is equal to

$$I_{sp}^{(4)} = i \sqrt{5} \left( \frac{3}{4} \right)^2, \quad \text{correspondingly } I_{sp}^{(4)} = i \sqrt{5} \left( \frac{3}{4} \right)^2 \frac{\tilde{Z} h}{a_0}.$$ (35)

In calculation of a transition probability in the unit time we will take into account, that at least three electrons participate in the matrix element of the operator (10) and the number of spin states in the determinant (3) is $2S + 1 = 4$:

$$d\omega_{if} = \frac{2 \pi}{h^2} (2S + 1) \left( \frac{e}{mc^2} \right)^2 \frac{2 \pi \hbar c^2}{\omega} |3\langle \Psi_{in} | \hat{p}_{1z} | \Psi_{Cr} \rangle|^2 \cos^2 \theta \delta(\omega_{if} - \omega) \frac{k^2 \delta k \delta \Omega}{(2\pi)^3}.$$ (36)

Here $\omega_{if}$ is appropriate to a transition frequency of a neutral radiation defect from the energy level of the corresponding free atom state on the energy level corresponding to a hybridized state in a lattice.

After integration on the wave vector of photons and averaging on all angles we get:

$$w_{if} = \begin{cases} 
(17/243) \tilde{Z}_B^3 \alpha^3 \omega_{if}^C & \text{for Diamond,} \\
(17 \cdot 5^3/2^3 3^6) \tilde{Z}_Si^3 \omega_{if}^Si & \text{for Silicon,} \\
(17 \cdot 15/212) \tilde{Z}_Ga^3 \omega_{if}^Ge & \text{for Germanium,} 
\end{cases}$$ (37)

where $\alpha = 1/137$ is the fine structure constant.
In accordance with Slater the effective charge is determined as \( \tilde{Z} = Z - \sigma \), where \( Z \) is the real nuclear charge and \( \sigma \) is a screening constant. For the atoms under consideration we have:

\[
\tilde{Z}_B \approx 2.6, \quad \tilde{Z}_{\text{Al}} \approx 3.5, \quad \tilde{Z}_{\text{Ga}} \approx 5.0.
\]  

(38)

4. Conclusion

We can see, that formulae (37) strongly differ from the well known “usual” formula for the radiation transition rate in atoms. The radiation transition rate (37) linearly depends on a transition frequency, square on an effective nuclear charge, but as usual contains \( \alpha^3 \).

Transition frequencies are unknown and must be calculated numerically, but for qualitative estimations we can take appropriate values for a forbidden gap in crystals. So, we have

\[
w_{if} = \begin{cases} 
5.5 \cdot 10^6 \text{e}^{-1} & \text{for Diamond}, \\
3.8 \cdot 10^8 \text{e}^{-1} & \text{for Silicon}, \\
4.9 \cdot 10^8 \text{e}^{-1} & \text{for Germanium}, 
\end{cases}
\]

(39)

Numerical values (39) show, that formation time of an acceptor center induced by negative muon in crystal is sufficiently long and must be taken into account in interpretation of all experimental results.

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References

[1] Smilga V P and Belousov Yu M 1994 The Muon Method in Science (New York: Nova Science)
[2] Vaisenberg A O 1964 Mu-meson (Moscow: Nauka) (in Russian)
[3] Imazato J, Nagamine K, Yamazaki T et al 1984 Phys. Rev. Lett. 53 1849
[4] Keller H, Kiefl R F, Baumeler H P et al 1986 Hyp. Int. 31 461
[5] Koch M, Meier K, Major J et al 1990 Hyp. Int. 65 1039
[6] Gorelkin V N, Grebinnik V G et al 1993 Nuclear Physics 56, 29
[7] Mamedov T N, Chaplygin I L, Duginov V N et al 1997 Hyp. Int. 105 345
[8] Mamedov T N, Gritsaj K I, Stoykov A V et al 2000 Physica B., 289-290 574
[9] Brewer J H, Krasnoperov E P, Kreitzman S R et al 1991 JETP Lett. 53 600
[10] Higemoto W, Sotoh K, Nishida N et al 1997 Hyp. Int. 106 39
[11] Krasnoperov E P, Meylikhov E E, Baines C et al 1996 Hyp. Interactions 97/98 347
[12] Belousov Yu M and Smilga V P 1994 JETP 79 811
[13] Belousov Yu M and Smilga V P 1999 Physica B 289-290 499
[14] Belousov Yu M 2007 JETP 131 243 (in Russian)
[15] Mamedov T N et al Dubna 2007 Study of Boron acceptor center in synthetic diamond by \( \mu \)SR-technique (Preprint JINR P14-2007-12) (in Russian)
[16] Baturin A S, Gorelkin V N, Solov’ev V R et al Dubna 2007 Particularity of formation of an acceptor center by negative muon in Diamond and its polarization behaviour (Preprint JINR) (in Russian)