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THE MUON CAPTURE BY ^{11}B TO THE TWO LOWEST LEVELS OF ^{11}Be

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A B S T R A C T

The partial rates for the capture of muons in ^{11}B for the transitions to the two lowest levels of ^{11}Be are calculated with the simple shell model. Although a comparison of the theoretical with the experimental result for the ground state to ground state beta decay shows that only qualitative agreement is expected, it turns out that the experiment of the muon capture can provide useful information on the first excited state of ^{11}Be , especially on its parity.

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1. INTRODUCTION

Talmi and Unna ¹⁾ have given some time ago a nice explanation of the experimental indication that the order of the two lowest levels in ^{11}Be is reversed with what one expects from the simple shell model description : the ground state should have the odd neutron in the $2s_{\frac{1}{2}}$ state instead of in the $1p_{\frac{1}{2}}$ state. The first excited state, at 320 KeV should then have the odd neutron in the $1p_{\frac{1}{2}}$ state. The experimental evidence is the very high $\log ft$ - value (≈ 6.7) for the beta decay of the ground state of ^{11}Be to the two lowest levels of ^{11}B which have the odd proton in the $1p_{\frac{3}{2}}$ and $1p_{\frac{1}{2}}$ state respectively ²⁾.

Recently Deutsch et al. suggested that experimental determination of the muon capture rate in ^{11}B to the two lowest states of ^{11}B should support the suggestion of Talmi and Unna. In particular this experiment should be able to confirm their prediction about the spin and parity of the first excited state.

We present in this paper the results of the calculation of the muon capture by ^{11}B to the ground state and the first excited state of ^{11}Be . We assume the harmonic oscillator shell model for the nuclear wave functions. Of course one does not expect that we obtain very accurate predictions with this simple model because it is just the interaction of the odd nucleon with the core that causes the order of the energy levels to be reversed. To find to what extent and in which way experimental results may be interpreted with this simple model, we calculate also the beta decay rate (first forbidden) of ^{11}Be (odd neutron in $2s_{\frac{1}{2}}$ state) to the ground state of ^{11}B (odd proton in $1p_{\frac{3}{2}}$ state)

The conclusion will be that a more sophisticated model for the nuclear wave functions is needed for a quantitative description.

In Section 2 we give the results for the muon capture. Section 3 contains the calculation of the beta decay of ^{11}Be . Section 4 gives a discussion of the results.

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2. THE MUON CAPTURE

The method of calculation is analogous to the treatment of the partial muon capture in ^{16}O in Ref. ³⁾, to which we refer also for an explanation of the notation.

The difference with the $^{16}\text{O} \rightarrow ^{16}\text{N}^{(*)}$ transitions is that we have there essentially one particle involved whereas in the case of $^{11}\text{B} \rightarrow ^{11}\text{Be}^{(*)}$ each of the three protons in the $1p_{3/2}$ shell of the ^{11}B ground state may make the transition to a neutron in the $1p_{1/2}$ or $2s_{1/2}$ shell in the final state, leaving the other two protons in the $1p_{3/2}$ shell. We assume that the wave functions for the nucleons in the filled (sub) shells are not changed in the transition.

We show in the Appendix that for our case the matrix element of a one-particle operator $\mathcal{O}$ between initial and final nuclear states ψ_i and ψ_f is given by

$$\langle \psi_f | \sum_{i=1}^A \mathcal{O} | \psi_i \rangle = \frac{1}{\sqrt{2}} \langle \psi_{J_f M_f} | \mathcal{O} | \psi_{J_i M_i} \rangle \quad (1)$$

where $\psi_{J_f M_f}$ and $\psi_{J_i M_i}$ are the one particle harmonic oscillator wave functions for the odd nucleon in final and initial state respectively.

The rate of the muon capture from a nuclear state a to a state b is given by, Ref. ³⁾

$$\Lambda_{\mu c}(a \rightarrow b) = \frac{r_{ab}^2}{2\pi} \left| \frac{1}{\mu} \right|_{av}^2 M_{ab}^2 \quad (2)$$

The squared matrix element M_{ab}^2 may be written as

$$M_{ab}^2 = \frac{1}{(2j_a + 1)} \sum_{M_a M_b} \int \frac{d\Omega_{\vec{v}}}{4\pi} \{A_0 + A_1 + A_2\} \quad (3)$$

where

$$A_0 = g_V^2 |\int 1|^2 + g_A^2 |\int \vec{\sigma}|^2 + (g_P^2 - 2g_P g_A) |\int \hat{v} \cdot \vec{\sigma}|^2 \quad (4)$$

$$A_1 = -g_V g_A \{ (\int 1)^* (\int \frac{\vec{p} \cdot \vec{v}}{M}) + c.c. \} - (g_A^2 - g_P^2) g_A \{ (\int \vec{v} \cdot \vec{\sigma})^* (\int \frac{\vec{p} \cdot \vec{\sigma}}{M}) + c.c. \} \\ + g_A g_V i \hat{v} \cdot \{ (\int \vec{\sigma})^* \wedge (\int \frac{\vec{p}}{M}) - c.c. \} \quad (5)$$

$$A_2 = g_A^2 |\int \frac{\vec{\sigma} \cdot \vec{p}}{M}|^2 + g_V^2 |\int \frac{\vec{p}}{M}|^2 \quad (6)$$

We have used the abbreviation

$$\int 0 = \langle j_b M_b | \sum_i \tau_i^{(-)} \exp(-i \vec{v}_{ab} \cdot \vec{r}_i) 0_i | j_a M_a \rangle \quad (7)$$

For the transition in which a $1p_{3/2}$ proton changes into a $2s_{1/2}$ neutron (ground state $\rightarrow$ ground state) we obtain

$$M_{1p_{3/2} \rightarrow 2s_{1/2}}^2 = \frac{1}{2} \{ g_V^2 + 2g_A^2 + (g_A - g_P)^2 \} \int_1^2 \\ - g_V g_A (\int_2 - \frac{2}{3} \int_3) \int_1 - (g_A - g_P) g_A \int_3 \int_1 + g_A g_V \int_2 \int_1 \quad (8)$$

For the transition to the first excited state with the odd neutron in the $1p_{\frac{1}{2}}$ state we have

$$\begin{aligned}
 M_{1p_{\frac{3}{2}} \rightarrow 1p_{\frac{1}{2}}}^2 = & \frac{1}{2} \int_A \left(\frac{4}{3} P_0^2 + \frac{5}{3} P_2^2 \right) + \frac{1}{2} \left(\int_\rho^2 - 2 \int_\rho \int_A \right) \left(\frac{4}{9} P_0^2 - \right. \\
 & \left. - \frac{4}{9} P_0 P_2 + \frac{5}{9} P_2^2 + \frac{1}{2} \int_V^2 P_2^2 - \int_V \int_V P_1 P_2 + \right. \\
 & \left. + \left(\int_A - \int_\rho \right) g_A \times \frac{2}{3} P_3 \left(P_0 - \frac{1}{4} P_2 \right) + \frac{2}{9} \int_A \int_V P_3 P_0 \right. \quad (9)
 \end{aligned}$$

Apart from the radial integrals I_1 , I_2 and I_3 for the $1p_{\frac{3}{2}} \rightarrow 2s_{\frac{1}{2}}$ transition, already defined in Ref.3), we introduced in (9) the following integrals for the $1p_{\frac{3}{2}} \rightarrow 1p_{\frac{1}{2}}$ transition :

$$P_0 = \int R_{1p}(r) f_0(r) R_{1p}(r) r^2 dr = \left(1 - \frac{1}{6} p^2\right) \exp\left(-\frac{1}{4} p^2\right)$$

$$P_1 = \frac{1}{5M} \int R_{1p}(r) \left(\frac{3}{r} + 2 \frac{\partial}{\partial r}\right) (R_{1p}(r)) r^2 dr = \frac{1}{30} p^3 g \exp\left(-\frac{1}{4} p^2\right)$$

$$P_2 = \int R_{1p}(r) f_2(r) R_{1p}(r) r^2 dr = \frac{1}{6} p^2 \exp\left(-\frac{1}{4} p^2\right)$$

$$P_3 = \frac{1}{M} \int R_{1p}(r) f_3(r) r \frac{\partial}{\partial r} \left(\frac{1}{r} R_{1p}(r)\right) r^2 dr = -\frac{5}{6} p g \left(1 - \frac{1}{10} p^2\right) \exp\left(-\frac{1}{4} p^2\right)$$

We omitted terms quadratic in the nucleon velocity (A_2) ; these have a negligible influence ($\leq 1\%$). In the numerical evaluation we have taken for the harmonic oscillator parameter b the value determined in the electron scattering experiments ⁴⁾

$$b = (1.55 \pm .15) \text{ fm} \quad (10)$$

The average of the muon wave function over the nuclear volume, $|\varphi_\mu|_{\text{av}}^2$ is obtained from an interpolation of the results of Sens for Be and C ⁵⁾

$$|\varphi_\mu|_{\text{av}}^2 = .92 \times \left(\frac{5\pi\mu e^2}{\pi} \right)^3 \quad (11)$$

The results for several interesting cases are summarized in the table. The discussion of these results will be given in Section 4.

3. THE BETA DECAY

We calculate in this Section the rate for the beta decay of the ground state of ^{11}Be (with the odd neutron in the $2s_{\frac{1}{2}}$ state) to the ground state of ^{11}B (with the odd proton in the $1p_{\frac{3}{2}}$ state). We compare the result with the experimental value. Then we are able to say something about the adequacy of the model calculations in Section 2, because the matrix elements occurring in the muon capture and beta decay are quite similar.

The calculation of the beta decay is somewhat more complicated than the calculation of the analogous muon capture transition, because we have to use the electron wave function in the Coulomb field of the nucleus. We use the results of Lee-Whiting ⁶⁾. The decay rate for a transition of the nucleus in state a to state b is given by (first-forbidden transition !)

$$\Lambda_{\beta}(a \rightarrow b) = \frac{1}{2\pi^3} \int_0^{E_M} dE \, p \, E (E_M - E)^2 F_0(Z, E) S_1(a \rightarrow b) \quad (12)$$

E_M is the maximum energy of the electron ($=11.5 \text{ MeV}$), E and p are the energy and momentum of the electron. $F_0(Z, E)$ is the Fermi function. Because Z is very small we use the approximation

$$F_0(Z, E) = \frac{2\gamma\pi}{1 - \exp(-2\gamma\pi)} \quad (13)$$

with

$$\gamma = \frac{\alpha Z E}{p}$$

Using formula (106) of Lee-Whiting⁶⁾ we obtain for the form-factor $S_1(a \rightarrow b)$ in the case of a $(\frac{1}{2}^+ \rightarrow \frac{3}{2}^-)$ transition

$$\begin{aligned}
 S_1(\frac{1}{2}^+ \rightarrow \frac{3}{2}^-) = & g_V^2 \{ | \langle 311 \rangle |^2 (M_0 - \frac{2}{3} q N_0 + \frac{1}{3} q^2 L_0 + 2L_1) + \\
 & + | \langle \alpha 10 \rangle |^2 L_0 + 2 \langle 311 \rangle \langle \alpha 10 \rangle^* (N_0 - \frac{1}{3} q L_0) \} \\
 & + g_A^2 \{ | \langle \sigma 11 \rangle |^2 (M_0 + \frac{2}{3} q N_0 + \frac{1}{3} q^2 L_0 + \frac{1}{2} L_1) + | \langle \sigma 21 \rangle |^2 (\frac{1}{3} q^2 L_0 + 3L_1) \} \\
 & + 2g_V g_A \{ \langle \sigma 11 \rangle \langle 311 \rangle^* (M_0 - L_1) + \langle \sigma 11 \rangle \langle \alpha 10 \rangle^* (N_0 + \frac{1}{3} q L_0) \} \quad (14)
 \end{aligned}$$

$q = E_M - E$ is the momentum of the neutrino.

The nuclear matrix elements $\langle 111 \rangle$, $\langle \alpha 10 \rangle$ etc. are proportional to the reduced matrix elements of the occurring operators. These reduced matrix elements are quite similar to the matrix elements which we had to compute for the muon capture. So we have for example

$$\langle \alpha 10 \rangle = -(\pi)^{\frac{1}{2}} \langle 1 p_{3/2} || T(1) || 2 s_{1/2} \rangle \quad (15)$$

where

$$T(1) = \frac{p_N}{M}$$

(p_N is nucleon momentum operator).

Thus this matrix element is proportional to the matrix element $\int \frac{p}{M}$ of Section 2, where we have to take only the first non-vanishing term in the expansion of $\exp(-i \vec{v} \cdot \vec{r})$.

M_0 , N_0 and L_0 , are combinations of the electron Coulomb wave function taken at the "nuclear radius" R . Greuling⁷⁾ gives the

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first terms of the expansion of these expressions in powers of R in the limit $(\alpha Z)^2 \ll 1$:

$$\begin{aligned} M_0 &= \beta^2 + \frac{2}{3} \beta \frac{p^2}{E} + \frac{1}{9} p^2 \\ N_0 &= -\beta - \frac{p^2}{3E} \\ L_0 &= 1 \\ L_1 &= \frac{1}{9} p^2 \end{aligned} \quad (16)$$

where

$$\beta = \frac{\alpha Z}{2R}$$

Expressing the matrix elements in terms of the harmonic oscillator parameter b we obtain

$$\begin{aligned} S_1\left(\frac{1}{2}^+ \rightarrow \frac{3}{2}^-\right) &= g_V^2 \left\{ \frac{1}{6} b^2 (M_0 - \frac{2}{3} g N_0 + \frac{1}{3} g^2 L_0 + 2 L_1 + \right. \\ &\quad \left. + \frac{1}{6(Mb)^2} + \frac{1}{3M} (N_0 - \frac{1}{3} g L_0) \right\} + \\ &\quad + g_A^2 \left\{ \frac{1}{6} b^2 (M_0 + \frac{2}{3} g N_0 + g^2 L_0 + 0 L_1) \right\} + \\ &\quad + \frac{1}{3} g_V g_A \left\{ b^2 (M_0 - L_1) + \frac{1}{M} (N_0 + \frac{1}{3} g L_0) \right\} \end{aligned} \quad (17)$$

We have taken for the "radius" of the nucleus appearing in β (16) the correct averaged value obtained with the harmonic oscillator wave functions

$$R = \frac{3}{2} \sqrt{\pi} b \quad (18)$$

The result for Λ_β becomes

$$\Lambda_\beta^{\text{theor}}\left(\frac{1}{2}^+ \rightarrow \frac{3}{2}^-\right) = (4.5 \pm 1.2) \times 10^{-2} s^{-1} \text{ for } b = (1.55 \pm 0.15) \text{ fm.} \quad (19)$$

If we neglect the influence of the Coulomb terms ($\beta = 0$) we obtain

$$\Lambda_{\beta} \left(\frac{1}{2}^+ \rightarrow \frac{3}{2}^- \right) = 4.4 \times 10^{-2} s^{-1} \text{ for } b = 1.55 \text{ fm} \quad (20)$$

This is almost equal to the result of Eq. (19) : however this appears to be an accidental cancellation of different relatively large terms.

4. DISCUSSION

Comparing the theoretical result for the beta decay (19) with the experimental value ²⁾

$$\Lambda_{\beta}^{\text{exp}}\left(\frac{1}{2}^{+} \rightarrow \frac{3}{2}^{-}\right) = 3.1 \times 10^{-2} \text{ s}^{-1} \quad (21)$$

obtained from the half-life of ^{11}Be and the branching ratio to the ground state of ^{11}B , we see that the theoretical value is about a factor 2.4 too large. This is not surprising because shell model calculations of transition matrix elements deviate generally quite a lot from the experimental results. The model used is too simple. Especially in our case this is evident because it is just the interaction of the $2s_{\frac{1}{2}}$ and $1p_{\frac{1}{2}}$ neutron with the other nucleons which causes the order of the two lowest levels in ^{11}Be to be reversed ¹⁾. To improve the theoretical result, this should be taken into account. However the calculation becomes much more complicated. Furthermore it occurs often that transition matrix elements remain very badly predicted although energy levels fit very well (see e.g. ³⁾). The reason may be that one does not know the right one-particle wave functions. Further the assumption that the nucleons in the closed (sub) shells may be neglected might be not justified. If, for example, the value of the harmonic oscillator parameter is different in the initial and final nucleus, the cores do not overlap and may cause a reduction of the matrix element.

Another point to be made is that for light nuclei the L-S coupling scheme works at least as well as the j-j scheme. The reason why we used j-j coupling is that the spin of the ground state of ^{11}B is predicted wrongly in L-S coupling ($J = \frac{1}{2}$, cf. e.g., ⁹⁾). Because of the above considerations it is clear that the comparison of the muon capture results with experiment will be rather ambiguous. We make the following remarks :

- i) We see from the table that the ratio of the rates for capture into the $1p_{\frac{3}{2}}$ and the $2s_{\frac{1}{2}}$ states is about 7. This shows that the distinction between allowed (transition to $1p_{\frac{1}{2}}$ state) and first-forbidden (transition to $2s_{\frac{1}{2}}$ state) in muon capture is less clear than in beta decay where we have generally about a factor 100 - 1000. The reason is that the neutrino momentum in muon capture (≈ 90 MeV/c) is much larger than the maximum electron momentum (in our case ≈ 10 MeV/c), so that the first terms of the expansion of the exponential differ less in muon capture compared with beta decay.
- ii) Although the theoretical and experimental results for the beta decay show that our model calculation cannot be trusted, one may nevertheless hope that the experimental results for the muon capture will show a difference of a factor 4 - 15 between the two capture rates. There is no reason to expect a more quantitative agreement because the final states are quite different in the two transitions. This shows itself also in the rather strong dependence on b of the ratio of the two capture rates.
- iii) The table shows that the ratio of the two rates does almost not depend on the pseudoscalar coupling constant whereas this dependence of each of the rates separately is not very spectacular and not different from more reliable cases as there are captures in ${}^3\text{He}$ and ${}^{12}\text{C}$.
- iv) One may hope that a comparison of the ratio of the rates for muon capture and beta decay between the ground states of ${}^{11}\text{B}$ and ${}^{11}\text{Be}$ will provide a more quantitative agreement between experiment and theory. However, although the initial and final nuclear states are the same, this ratio depends rather strongly on b and thus on the (poor) model used. The reason is that the velocity terms are only small corrections in the muon capture whereas they form the major contribution in the beta decay. Nevertheless some cancelling, e.g., because of non-overlapping cores, may occur.

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A P P E N D I X

We derive expression (1) for the matrix element of a one-particle operator. The effective part of the initial state has three protons in the $1p_{\frac{3}{2}}$ shell which couple to a total spin $J_1 = \frac{3}{2}$. Using formula (26.10) of de-Shalit and Talmi⁽⁸⁾, in the following abbreviated by ST, we have

$$\psi(j_i, M_i) = \sum_{j_1, \text{even}} [f^{a_2}(j_1) f^a j_i | f^{a_3} j_i] \psi(f^{a_2}(j_1) f^{a_3} j_i, M_i) \quad (\text{A.1})$$

where [ST(25.1)]

$$\psi(f^{a_2}(j_1) f^{a_3} j_i, M_i) = \sum_{m_1, m_2, m_3} \langle f^{a_2} m_1, f^{a_3} m_2 | j_1, M_1 \rangle \times$$

$$\times \langle j_1, M_1, f^{a_3} m_3 | j_i, M_i \rangle \psi_{f^{a_2} m_1}(x_1) \psi_{f^{a_3} m_2}(x_2) \psi_{f^{a_3} m_3}(x_3) \quad (\text{A.2})$$

$\psi_{jm}(x)$ are the one-particle harmonic oscillator wave functions in our case for the $1p_{\frac{3}{2}}$ shell, $j^a = \frac{3}{2}$.

The coefficient of fractional parentage with $J_1 = 0$ is given by [ST (26.11) a]

$$[f^{a_2}(0) f^{a_3} \frac{3}{2} | f^{a_3} \frac{3}{2}] = \sqrt{\frac{1}{6}} \quad (\text{A.3})$$

whereas we do not need this coefficient for $J_1 = 2$. The final state has two protons in the $1p_{\frac{3}{2}}$ shell and one neutron in the $2s_{\frac{1}{2}}$ or $1p_{\frac{1}{2}}$ shell.

The only way to get a final spin $J_f = \frac{1}{2}$ is to couple the two protons to spin 0 which is then coupled with the spin of the neutron. Equation [ST(26.3)] gives with the appropriate normalization factor

$$\psi_f(j_f, M_f) = \frac{1}{\sqrt{3}} \left[\psi(j_{12}^a(0) j_3^c j_f M_f) - \psi(j_{13}^a(0) j_2^c j_f M_f) + \psi(j_{23}^a(0) j_1^c j_f M_f) \right] \quad (\text{A.4})$$

where, for our case [cf. ST(25.1)]

$$\begin{aligned} \psi(j_{12}^a(0) j_3^c j_f M_f) &= \sum_{m_a} (j_a^{m_a} j^{a-m_a} 100) \times \\ &\times (00 j_c^{m_c} j_f M_f) \psi_{j_a^{m_a}}(x_1) \psi_{j^{a-m_a}}(x_2) \psi_{j_c^{m_c}}(x_3) = \\ &= - \sum_{m_a} (j_a^{m_a} j^{a-m_a} 100) \psi_{j_a^{m_a}}(x_1) \psi_{j^{a-m_a}}(x_2) \psi_{j_c^{m_c}}(x_3) \end{aligned} \quad (\text{A.5})$$

and $j^a = \frac{3}{2}$ and $j^c = \frac{1}{2}$.

We now calculate the matrix element

$$\langle \psi_f(j_f M_f) | \sum_{i=1}^A Q_i | \psi_i(j_i M_i) \rangle =$$

$$3 \langle \psi_f(j_f M_f) | Q_3 | \psi_i(j_i M_i) \rangle =$$

$$\frac{3}{\sqrt{3}} \langle \psi(j_{12}^a(0) j_3^c j_f M_f) | Q_3 | \psi_i(j_i M_i) \rangle \quad (\text{A.6})$$

The last equality follows from the fact that only the term in ψ_f (A.4) which contains $\psi_{j c M_f}(\mathbf{x}_3)$ contributes. Because of the orthonormality of the oscillator wave functions we have

$$m_1 = m_a, m_2 = -m_a, M_1 = 0 \text{ and } m_3 = M_i.$$

The contribution to (A.6) of the term in ψ_i [cf. (A.1)] with $J_1 = 0$ becomes

$$\frac{3}{13} \times \frac{1}{\sqrt{6}} \sum_{m_a} (f^{a m_a} f^{a - m_a} | 0 0) (f^{a m_a} f^{a - m_a} | 0 0) \times$$

$$\times \langle \psi_{j c M_f} | 0 | \psi_{j a M_i} \rangle = \frac{1}{\sqrt{2}} \langle \psi_{j c M_f} | 0 | \psi_{j a M_i} \rangle \quad (\text{A.7})$$

The term with $J_1 = 2$ does not contribute because

$$\sum_{m_a} (f^{a m_a} f^{a - m_a} | 0 0) (f^{a m_a} f^{a - m_a} | 2 0) = 0$$

Thus, our final result is

$$\langle \psi_{j c M_f} | \sum_i \sigma_i | \psi_{j a M_i} \rangle = \frac{1}{\sqrt{2}} \langle \psi_{j c M_f} | 0 | \psi_{j a M_i} \rangle \quad (\text{A.8})$$

We see that we get a factor $\frac{1}{\sqrt{2}}$ compared with the too simple version of the shell model which allows only the odd proton to make the transition.

	$b(\text{fm})$	$\frac{g_P}{g_A}$	$\Lambda_{\mu c}(\frac{3}{2}^- \rightarrow \frac{1}{2}^+) \times 10^{-2} \text{ s}^{-1}$	$\Lambda_{\mu c}(\frac{3}{2}^- \rightarrow \frac{1}{2}^-) \times 10^{-2} \text{ s}^{-1}$	$\Lambda_{\mu c}(\frac{3}{2}^- \rightarrow \frac{1}{2}^-) / \Lambda_{\mu c}(\frac{3}{2}^- \rightarrow \frac{1}{2}^+)$
Without velocity terms	1.55	7	0.46	3.21	7.0
	1.55	0	0.62	4.05	6.5
Velocity terms included		7	0.54	3.44	6.3
		14	0.49	3.06	6.2
	1.40	7	0.51	3.67	7.2
	1.70	7	0.56	3.23	5.8

- Table -

The muon capture by ^{11}B into the two lowest levels of ^{11}Be for different values of the ratio g_P/g_A and the harmonic oscillator parameter b .

R E F E R E N C E S

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The beta decay of ^{11}Be to ^{11}B :
 between parentheses the (probable)
 state of the odd nucleon