DEPOLARIZATION OF THE NEGATIVE MUON IN MESIC-ATOM TRANSITIONS

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A method is developed for calculating the depolarization of μ^--mesons in mesic-atom cascade transitions. An estimate of the expected μ^--meson depolarization is given, which is in good agreement with available experimental data.

As noted by Gol'dman,[1] a μ^--meson is not substantially depolarized in Coulomb collisions. The spin-orbit interaction in mesic atom transitions[2] is mainly responsible for the depolarization.* This results from the fact that for a mesic atom level with fine-structure splitting ν and level width 2γ, the ratio always satisfies the inequality

$$(\nu/2\gamma)^2 \gg 1, \quad (1)$$

i.e., the period of precession is much less than the lifetime of the level.

Because of its large mass, in comparison with that of the electron, a slow μ^- meson is captured into a highly-excited quantum state of the atom. The principal quantum number n of such a state can be determined[3] from the approximate equality of Bohr radii of the meson and electron, giving $n \sim \sqrt{\mu} \sim 15$.

As noted by Bohr,[4] for a given quantum number n, the meson will be captured in a circular orbit ($l = n - 1$) with greatest probability, because of the large statistical weight of this orbit. The comparison with experimental data carried out by Stearns[5] and the subsequent remarks of Day and Morrison[6] indicate that both π^- and μ^- mesons are mainly captured in orbits with high l.

Following this, as a result of subsequent Auger transitions and radiative transitions, the μ^- meson loses energy and drops to the 1s level, from which it either is captured by the atom or decays, depending on Z.

Burbidge and De Borde[3,7] calculated the probability of mesic-atom transitions. According to their calculations, the picture is the following. For large n, dipole Auger transitions with $\Delta n = \Delta l = -1$ take place. Following this, for smaller n (if Z is not very small), the electric-dipole radiative transitions are essential, in which the circular orbits become even more probable because of the maximum probability of a radiative transition to a level with smallest possible n. Thus, in both Auger and radiative transitions, electric-dipole transitions are essential, in which the orbital quantum number decreases by unity ($\Delta l = -1$).

Starting from the above, we consider the depolarization of a μ^- meson in electric dipole transitions. Let the μ^- meson go from a level l_N to a level l_1 by successive dipole emission of either γ-rays or Auger electrons; that is, a cascade down $l_N(1)\, l_{N-1}(1)\, l_{N-2}(1) \ldots l_1(1)$ takes place. To obtain the corresponding density matrix, it is necessary to solve the equations by means of the perturbation theory given by Wigner and Weisskopf,[8] which takes into account the finite width of the level. However, as Abragam and Pound[9] have shown for the case of a single level, one can obtain the final result from the following simple considerations.

The interaction can be included in the expression for the transition probability by means of the operator $\mathbf{u} = \exp(-i\mathbf{K}t/\hbar)$ which operates on the wave functions of the meson in the given level ($\mathbf{K}$ here is the spin-orbit interaction energy operator). It is then necessary to consider the possibility of radiation in the interval t, t + dt, with probability $\exp(-2\gamma t)\, 2\gamma dt$. Integrating such a "supplemented" transition probability over t from 0 to ∞, we obtain the factor $[1 + (\nu_{jj'}/2\gamma)^2]^{-1}$ for each level, leading to the following density matrix

*I. M. Shmushekvich called my attention to this.

$$\rho = \sum a_{\sigma_N} \frac{(j_{N+1} l_{N+1} \mu_{N+1} \sigma_N | H_{N+1} | j_N l_N \mu_N \sigma_N)(j_{N+1} l_{N+1} \mu_{N+1} \sigma_N | H_{N+1} | j'_N l_N \mu'_N \sigma_N)^*}{1 + (\nu_{j_N j'_N}/2\gamma_N)^2}$$
$$\times \frac{(j_N l_N \mu_N \sigma_{N-1} | H_N | j_{N-1} l_{N-1} \mu_{N-1} \sigma_{N-1})(j'_N l_N \mu'_N \sigma'_{N-1} | H_N | j'_{N-1} l_{N-1} \mu'_{N-1} \sigma'_{N-1})^*}{1 + (\nu_{j_{N-1} j'_{N-1}}/2\gamma_{N-1})^2}$$
$$\ldots \times \frac{(j_2 l_2 \mu_2 \sigma_1 | H_2 | j_1 l_1 \mu_1 \sigma_1)(j'_2 l_2 \mu'_2 \sigma'_1 | H_2 | j'_1 l_1 \mu'_1 \sigma'_1)^*}{1 + (\nu_{j_1 j'_1}/2\gamma_1)^2} \qquad (2)$$
$$\times (j_1 l_1 \mu_1 \sigma | H_1 | j l \mu \sigma)(j'_1 l_1 \mu'_1 \sigma | H_1 | j l \mu \sigma)^*.$$

The summation here is over all unobserved properties of the radiation and mesic atom, and also over the two values of total angular momentum j_i, j'_i, over the projections of the total angular momentum μ_i, μ'_i and spin projections σ_i, σ'_i of the μ^- meson for each of the i levels ($i = 1, \ldots, N$); a_σ is the probability of the value σ_N for the spin projection on the symmetry axis before the cascade. The levels l_{N+1} and l are introduced formally in order to describe the formation and annihilation of the mesic atom (the decay of the μ^- meson or its capture by the atom).

Because of the inequality (1), all interference terms drop out of the Eq. (2) for the density matrix, and for each transition $l_{i+1}(1)\, l_i$, one can write[9]

$$(j_{i+1} l_{i+1} \mu_{i+1} \sigma_i | H_{i+1} | j_i l_i \mu_i \sigma_i)(j_{i+1} l_{i+1} \mu'_{i+1} \sigma'_i | H_{i+1} | j_i l_i \mu'_i \sigma'_i)^*$$
$$= (l_{i+1} | x | l_i)^2 \sum_{m_{i+1}, m_i, m'_{i+1}, m'_i} (-1)^{2\mu_{i+1} + 2\mu_i + m_{i+1} + m'_{i+1}}$$
$$\times (2l_{i+1} + 1)(2j_{i+1} + 1)(2j_i + 1) \qquad (3)$$
$$\times \begin{pmatrix} l_{i+1} & 1/2 & j_{i+1} \\ m_{i+1} & \sigma_i & -\mu_{i+1} \end{pmatrix} \begin{pmatrix} l_i & 1 & l_{i+1} \\ m_i & M_{i+1} & -m_{i+1} \end{pmatrix}$$
$$\times \begin{pmatrix} l_i & 1/2 & j_i \\ m_i & \sigma_i & -\mu_i \end{pmatrix} \begin{pmatrix} l_{i+1} & 1/2 & j_{i+1} \\ m'_{i+1} & \sigma'_i & -\mu_{i+1} \end{pmatrix}$$
$$\times \begin{pmatrix} l_i & 1 & l_{i+1} \\ m'_i & M_{i+1} & -m'_{i+1} \end{pmatrix} \begin{pmatrix} l_i & 1/2 & j_i \\ m'_i & \sigma'_i & -\mu_i \end{pmatrix} \delta_{\mu_{i+1}, \mu'_{i+1}} \delta_{\mu_i, \mu'_i}.$$

The round brackets denote Wigner 3j symbols. The $(l_{i+1} | x | l_i)$ denote factors which depend on the properties of the radiation and of the levels $l_{i+1} l_i$, but not on the magnetic quantum numbers. In the transition $l_{N+1} \rightarrow l_N$, carrying out summation in Eq. (3) over μ_{N+1}, M_{N+1} and j_{N+1}, we obtain an expression proportional to $(2j_N + 1) \times \begin{pmatrix} l_N & \frac{1}{2} & j_N \\ m_N & \sigma_N & -\mu_N \end{pmatrix}$; analogously, the transition $l_1 \rightarrow l$ gives the factor $(2j_1 + 1) \begin{pmatrix} l_1 & \frac{1}{2} & j_1 \\ m_1 & \sigma & -\mu_1 \end{pmatrix}^2$.

In the remaining cases, using the Racah relationship

$$\sum_{m_i, m_{i+1}, \sigma_i} (-1)^{m_{i+1}} \begin{pmatrix} l_i & 1/2 & j_i \\ m_i & \sigma_i & -\mu_i \end{pmatrix}$$
$$\times \begin{pmatrix} l_{i+1} & 1/2 & j_{i+1} \\ m_{i+1} & \sigma_i & -\mu_{i+1} \end{pmatrix} \begin{pmatrix} l_i & 1 & l_{i+1} \\ m_i & M_{i+1} & -m_{i+1} \end{pmatrix} \qquad (4)$$
$$= (-1)^{1/2 + 1 - \mu_i} \begin{Bmatrix} j_i & 1/2 & l_i \\ l_{i+1} & 1 & j_{i+1} \end{Bmatrix} \begin{pmatrix} j_i & 1 & j_{i+1} \\ \mu_i & M_{i+1} & -\mu_{i+1} \end{pmatrix},$$

where the curly brackets denote the Wigner 6j symbol, we obtain

$$\sum_{\sigma_i, \sigma'_i} (j_{i+1} l_{i+1} \mu_{i+1} \sigma_i | H_{i+1} | j_i l_i \mu_i \sigma_i)(j_{i+1} l_{i+1} \mu_{i+1} \sigma'_i | H_{i+1} | j_i l_i \mu_i \sigma'_i)$$
$$= (l_{i+1} | x | l_i)^2 (2l_{i+1} + 1)(2j_{i+1} + 1)(2j_i + 1) \qquad (5)$$
$$\times \begin{Bmatrix} j_i & 1/2 & l_i \\ l_{i+1} & 1 & j_{i+1} \end{Bmatrix}^2 \begin{pmatrix} j_i & 1 & j_{i+1} \\ \mu_i & M_{i+1} & -\mu_{i+1} \end{pmatrix}^2.$$

The mean value of the projection of the μ^- spin after the cascade of mesic-atom transitions, is determined by

$$\langle \sigma \rangle = \sum_\sigma \rho \sigma / \sum_\sigma \rho. \qquad (6)$$

We note that in Eq. (6) the factors corresponding to formation and annihilation of the mesic atom cancel, making it possible for us to formally introduce the levels l_{N+1}, l (see above).

Expressing σ in terms of the Wigner 3j symbol

$$\sigma = (-1)^{-1/2 + \sigma} \sqrt{\frac{3}{2}} \begin{pmatrix} 1/2 & 1 & 1/2 \\ \sigma & 0 & -\sigma \end{pmatrix} \qquad (7)$$

and then applying the Racah relationship, we have

$$\sum_{m_1, \sigma} (2j_1 + 1) \begin{pmatrix} l_1 & 1/2 & j_1 \\ m_1 & \sigma & -\mu_1 \end{pmatrix}^2 (-1)^{-1/2 + \sigma} \sqrt{\frac{3}{2}} \begin{pmatrix} 1/2 & 1 & 1/2 \\ \sigma & 0 & -\sigma \end{pmatrix} \qquad (8)$$
$$= (-1)^{-1/2 + l_1 + 1 - \mu_1} \sqrt{\frac{3}{2}} (2j_1 + 1) \begin{Bmatrix} j_1 & l_1 & 1/2 \\ 1/2 & 1 & j_1 \end{Bmatrix} \begin{pmatrix} j_1 & 1 & j_1 \\ \mu_1 & 0 & -\mu_1 \end{pmatrix}.$$

Employing Eq. (8) for each of the i levels ($i = 1, \ldots, N - 1$), we obtain

l_N	$(P/P_0)_{l_N, l_{N-1},\ldots 0}$	l_N	$(P/P_0)_{l_N, l_{N-1},\ldots,0}$
0	1	9	0.192
1	0.407	10	0.188
2	0.301	11	0.185
3	0.257	12	0.183
4	0.234	13	0.181
5	0.219	14	0.179
6	0.209	...	...
7	0.202	∞	0.156
8	0.197		

$$(2l_{i+1}+1)(2j_{i+1}+1)(2j_i+1)\begin{Bmatrix} j_i & 1/2 & l_i \\ l_{i+1} & 1 & j_{i+1}\end{Bmatrix}^2 \sum_{\mu_i, M_{i+1}} (-1)^{\mu_i} \times \begin{pmatrix} j_i & 1 & j_{i+1} \\ \mu_i & M_{i+1} & -\mu_{i+1}\end{pmatrix}^2 \begin{pmatrix} j_i & 1 & j_i \\ \mu_i & 0 & -\mu_i\end{pmatrix}$$
$$= (-1)^{1+1-\mu_{i+1}}(2l_{i+1}+1)(2j_{i+1}+1)(2j_i+1)\begin{Bmatrix} j_i & 1/2 & l_i \\ l_{i+1} & 1 & j_{i+1}\end{Bmatrix}^2 \times \begin{Bmatrix} j_{i+1} & 1 & j_i \\ j_i & 1 & j_{i+1}\end{Bmatrix}\begin{pmatrix} j_{i+1} & 1 & j_{i+1} \\ \mu_{i+1} & 0 & -\mu_{i+1}\end{pmatrix}. \quad (9)$$

Finally

$$(2j_N+1)\sum_{m_N, \mu_N}(-1)^{\mu_N}\begin{pmatrix} l_N & 1/2 & j_N \\ m_N & \sigma_N & -\mu_N\end{pmatrix}^2\begin{pmatrix} j_N & 1 & j_N \\ \mu_N & 0 & -\mu_N\end{pmatrix} \quad (10)$$
$$= (-1)^{l_N+1-\sigma_N}(2j_N+1)\begin{Bmatrix} 1/2 & l_N & j_N \\ j_N & 1 & 1/2\end{Bmatrix}\begin{pmatrix} 1/2 & 1 & 1/2 \\ \sigma_N & 0 & -\sigma_N\end{pmatrix}.$$

From the unitarity of the matrices of the Wigner 3j symbols, it is easy to show

$$\sum_\sigma \rho = \prod_i (l_{i+1}|x|l_i)^2(2l_N+1). \quad (11)$$

Thus, substituting Eqs. (2), (5), and (8) to (11) into Eq. (6) and using Eq. (7) for σ_N, we obtain

$$\frac{\langle\sigma\rangle}{\langle\sigma_N\rangle} \equiv \left(\frac{P}{P_0}\right)_{l_N, l_{N-1},\ldots, l_1} = \sum_{j_N, j_{N-1},\ldots, j_1}(-1)^{l_1+l_N+N-1}(2j_1+1)$$
$$\times\begin{Bmatrix} j_1 & l_1 & 1/2 \\ 1/2 & 1 & j_1\end{Bmatrix}(2l_2+1)(2j_2+1)(2j_1+1)\begin{Bmatrix} j_1 & 1/2 & l_1 \\ l_2 & 1 & j_2\end{Bmatrix}^2\begin{Bmatrix} j_2 & 1 & j_1 \\ j_1 & 1 & j_2\end{Bmatrix}$$
$$\times(2l_3+1)(2j_3+1)(2j_2+1) \quad (12)$$
$$\times\begin{Bmatrix} j_2 & 1/2 & l_2 \\ l_3 & 1 & j_3\end{Bmatrix}^2\begin{Bmatrix} j_3 & 1 & j_2 \\ j_2 & 1 & j_3\end{Bmatrix}\cdots(2j_N+1)(2j_{N-1}+1)$$
$$\times\begin{Bmatrix} j_{N-1} & 1/2 & l_{N-1} \\ l_N & 1 & j_N\end{Bmatrix}^2\begin{Bmatrix} j_N & 1 & j_{N-1} \\ j_{N-1} & 1 & j_N\end{Bmatrix}(2j_N+1)\begin{Bmatrix} 1/2 & l_N & j_N \\ j_N & 1 & 1/2\end{Bmatrix}.$$

We denote by P_0 and P the degree of polarization of the μ^- meson immediately before and after the cascade considered. Equation (12) gives the final polarization from which it is possible to find the depolarization, equal to $(1-P/P_0)\times 100\%$. In the future, we will, for simplicity, call P/P_0 the depolarization instead of $(1-P/P_0)$. We calculate P/P_0 for the cascade down $l_i = l_1 + i - 1$ $(i = 1, \ldots, N)$.

Rewriting the sum in Eq. (12) in the form

$$\left(\frac{P}{P_0}\right)_{l_N, l_{N-1},\ldots, l_1} = \sum_{\pm}\left(\frac{P}{P_0}\right)_{\underset{\pm}{l_N}, \underset{\pm}{l_{N-1}},\ldots, \underset{\pm}{l_1}}, \quad (13)$$

where the $\pm$ under l_i correspond to terms with $j_i = l_i \pm \frac{1}{2}$, it is easy to see that the only terms remaining in Eq. (13) are those in which all pluses stand to the right of the minuses. (Dipole transitions are possible only in this case.)

On the other hand, substituting in values for the Wigner 6j-symbols, we see that addition of pluses to the right of minuses to the left does not change the form of the corresponding term, so that there are terms of three types, which are easy to find:

$$\left(\frac{P}{P_0}\right)_{\underset{+}{l_N},\ldots, \underset{+}{l_i}, \underset{+}{l_{i-1}},\ldots, \underset{+}{l_1}} = \frac{(2l_N+3)(l_N+1)}{3(2l_N+1)^2};$$
$$\left(\frac{P}{P_0}\right)_{\underset{-}{l_N},\ldots, \underset{-}{l_i}, \underset{+}{l_{i-1}},\ldots, \underset{+}{l_1}} = -\frac{4l_i^2-5}{3(4l_i^2-1)^2}; \quad (14)$$
$$\left(\frac{P}{P_0}\right)_{\underset{-}{l_N},\ldots, \underset{-}{l_i}, \underset{-}{l_{i-1}},\ldots, \underset{-}{l_1}} = \frac{(2l_1-1)l_1}{3(2l_1+1)^2}.$$

Thus

$$\left(\frac{P}{P_0}\right)_{l_N, l_{N-1},\ldots, l_1} = \frac{1}{3}\left[\frac{(2l_N+3)(l_N+1)}{(2l_N+1)^2} - \sum_{l=l_1+1}^{l_N}\frac{4l^2-5}{(4l^2-1)^2} + \frac{(2l_1-1)l_1}{(2l_1+1)^2}\right]. \quad (15)$$

Setting $l_1 = 0$, we obtain the depolarization for various l_N from Eq. (15) (see table). As one would expect, the main contribution comes from low l. For example $(P/P_0)_{14,13,\ldots,2} = 0.237$.

Analogously, we can introduce the depolarization for other cascades:

$$\left(\frac{P}{P_0}\right)_{l_{N-1}-1, l_{N-1}, l_{N-1}-1, l_{N-1}-2,\ldots, l_1} = \frac{1}{3}\left\{\frac{(2l_{N-1}+3)(l_{N-1}+1)}{(2l_{N-1}+1)^2}\right.$$
$$+\frac{(2l_{N-1}+1)^2(2l_{N-1}-3)(l_{N-1}-1)-4l_{N-1}^2+5}{(2l_{N-1}-1)^3 l_{N-1}} \quad (16)$$
$$\left.\times\left[-\sum_{l=l_1+1}^{l_{N-1}}\frac{4l^2-5}{(4l^2-1)^2}+\frac{(2l_1-1)l_1}{(2l_1+1)^2}\right]\right\}.$$

From Eq. (16) it can be seen that the transition with $\Delta l = +1$ leaves the depolarization practically unchanged.

Thus for $l_{N-1} \gtrsim 5$, with an accuracy to the third decimal place,

$$(P/P_0)_{l_{N-1}-1, l_{N-1}, l_{N-1}-1, l_{N-1}-2,\ldots,0} = (P/P_0)_{l_{N-1}, l_{N-1}-1,\ldots,0},$$

If the transition with $\Delta l = +1$ takes place at the end of the cascade ($l_3 = l_1$, $l_2 = l_1 + 1$), then

$$\left(\frac{P}{P_0}\right)_{l_N, l_{N-1}, \ldots, l_1, l_1+1, l_1} = \frac{1}{3}\left\{\frac{1}{(2l_2+1)\,l_2}\left[\frac{(2l_2-1)^2(2l_2+3)(l_2+1)}{(2l_2+1)^2} + \frac{(4l_2^2-5)^2}{(4l_2^2-1)^2(2l_2-1)\,l_2} - \frac{(4l_2^2-5)(2l_2-3)(l_2-1)}{(2l_2-1)^3\,l_2}\right] \times \left[\frac{(2l_N+3)(l_N+1)}{(2l_N+1)^2} - \sum_{l=l_3+1}^{l_N}\frac{4l^2-5}{(4l^2-1)^2}\right] + \frac{(2l_2+1)^2(2l_2-3)(l_2-1)}{(2l_2-1)^3 l_2}\left[\frac{(2l_1-1)\,l_1}{(2l_1+1)^2} - \frac{4l_2^2-5}{(4l_2^2-1)^2}\right]\right\}. \tag{17}$$

In particular

$$(P/P_0)_{l_N, l_{N-1}, \ldots, 1, 0, 1, 0} = 0.407\,(P/P_0)_{l_N, l_{N-1}, \ldots, 0};$$
$$(P/P_0)_{l_N, l_{N-1}, \ldots, 1, 2, 1, 0} = 0.745\,(P/P_0)_{l_N, l_{N-1}, \ldots, 0} - 0.0178. \tag{18}$$

Substituting $l_N = 12$, we obtain

$$(P/P_0)_{12, 11, 10, \ldots, 0, 1, 0} = 0.074,$$
$$(P/P_0)_{12, 11, 10, \ldots, 1, 2, 1, 0} = 0.118.$$

respectively.

Calculation shows that if a quadrupole transition takes place for large l, the result does not differ from that given by Eq. (15).

Thus, the depolarization depends only weakly on the orbital moment of the initial level of the cascade and on the type of transition taking place for large l. This is understandable, in so far as in this case the deflection of $\mathbf{l}$ (and correspondingly of $\mathbf{s}$) which, roughly speaking, is proportional to the moment of force trying to turn $\mathbf{l}$ to l, is less for large l.

From the above, we can estimate the expected depolarization:

$$P/P_0 = \sum_i a_i (P/P_0)_i, \quad \sum_i a_i = 1, \tag{19}$$

where $(P/P_0)_i$ is the depolarization in the i-th cascade and a_i is the probability of this cascade. In Eq. (19) the sum is taken over all possible cascades. As already noted, the value of (19) is determined by the contribution of cascades beginning with large values of l_N.

In capture into a level $n = 15$, $l_N = 14$, a cascade with $\Delta n = \Delta l = -1$ proceeds, and the corresponding depolarization is, according to Eq. (15), equal to 0.179.

For $n = 15$, $l_N = 13$, Auger transitions with $\Delta n = \Delta l = -1$ will take place down to $n \sim 6$; following this, in the first radiative transitions, a transition into a circular orbit with $\Delta n = -2$, $\Delta l = -1$, will be most probable.[3]

Thus, in this case the depolarization will also be described by Eq. (15) in practice, and will be equal to 0.181. Ten times less probable is the cascade with $\Delta n = \Delta l = -1$ down to the 2s state, then[10] an Auger transition to the 2p state, followed by a radiative transition to the 1s state. In this cascade $(P/P_0)_i = 0.074$, i.e., in capture from the level $n = 15$, $l = 13$, $\langle P/P_0 \rangle = 18\%$.

For $n = 15$, $l_N = 12$, and also for smaller values of l, the cascade with $\Delta l = -1$ is again most important because of the probable transition to a circular orbit in the first radiative transitions.

As follows from Eq. (16), the possible transitions with $\Delta l = +1$ do not change matters at all, except for very small l. For small l, according to the Table, the value $(P/P_0)_{l_N, l_{N-1}, \ldots, 0}$ increases somewhat. However, the probability of transitions with $\Delta l = +1$ grows for small l, so that this increase will be compensated for. Thus, the expected value of the expression (18) should be $P/P_0 \sim 18\%$. This value agrees well with available experimental data[11] for carbon.

As seen from Eqs. (15) to (17), the sign of the polarization does not change in mesic atom transitions; these can thus be used to determine the decay scheme of the π^- meson.[2]

The influence of the electron shells in mesic atom transitions can be neglected, since within the lifetime of the levels of the mesic atom, the μ^- meson is not depolarized in this field. This is not true of the final 1s level, the lifetime of which is determined by the decay or capture of the μ^- meson. Taking into account this additional depolarization in the field of the electron shell of total angular momentum j_e, gives a factor $F = (4j_e^2 + 4j_e + 3)/3(2j_e + 1)^2$ in the right-hand sides of Eqs. (15) to (17). For example, in the case of carbon, the electron shells will be the configuration $1s^2 2s^2 2p$ with $j_e = \frac{1}{2}$, $F = \frac{1}{2}$, which gives $P/P_0 \sim 9\%$. However, in condensed material the influence of the neighboring atoms and electrons probably leads to a compensation of the magnetic field of the electron shells, i.e., to $j_e = 0$ and, consequently, to $F = 1$.

Equation (12) is applicable to a nucleus with zero spin. Taking into account the hyperfine structure leads to larger depolarization of the μ^- meson.

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39