Magnetic Field Dependence of Muonium-antimuonium Conversion

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We study the magnetic field dependence of muonium–antimuonium conversion induced by neutral (pseudo)scalar bosons. Only the $SS$ operator contributes to the conversion of polarized muonium, but it gets quenched by a magnetic field of strength 0.1 Gauss or stronger. Conversion induced by $SS$ couplings for unpolarized muonium is independent of magnetic field. Magnetic fields of 0.1 Tesla or stronger starts to suppress conversion induced by $PP$ interactions in the lowest Breit-Rabi level, but gets partially compensated by a rise in conversion probability in the other unpolarized level. The effects of $(S \mp P)(S \mp P)$ and $(S \mp P)(S \pm P)$ operators behave in the same way as $(V \mp A)(V \mp A)$ and $(V \mp A)(V \pm A)$ operators, respectively.
The spontaneous conversion of muonium (hydrogen-like atom $M = \mu^+e^-$) into antimuonium ($\bar{M}$) would violate the separate additive muon and electron numbers, but would remain consistent with multiplicative muon or electron number conservation. Defining the effective coupling $G_{M\bar{M}}$ via the interaction [1],

$$\mathcal{H}_{M\bar{M}} = \frac{G_{M\bar{M}}}{\sqrt{2}} \bar{\mu}\gamma\lambda(1 - \gamma_5)e\bar{\mu}\gamma^\lambda(1 - \gamma_5)e + \text{h.c.},$$  

the limit has just been improved [2] by an order of magnitude,

$$G_{M\bar{M}} < 1.8 \times 10^{-2} G_F,$$

compared with the previous bound of $0.16 G_F$ [3], where $G_F$ is the Fermi constant. The ultimate aim [4] is to reach the sensitivity level of $10^{-3} G_F$.

The $(V-A)(V-A)$ interaction of eq. (1) gained further theoretical footing when Halprin [5,6] pointed out that in left-right symmetric models with Higgs triplets, doubly charged scalars can mediate $M-\bar{M}$ transitions at tree level. The effective interaction is of $(V \pm A)(V \pm A)$ form after Fierz rearrangement. The possibility of $(V-A)(V+A)$ interactions, induced by dilepton gauge bosons [8], was discussed by Fuji et al. [7]. Interestingly, muonium conversion is more pronounced in the singlet channel, in contrast to the $(V-A)(V-A)$ case where the conversion matrix element is of equal strength for both singlet and triplet muonium [1].

Recently, Horikawa and Sasaki [9] pointed out further that $(V-A)(V+A)$ interactions are considerably less sensitive to magnetic fields compared to the $(V-A)(V-A)$ case. Since actual experiments are conducted in the presence of magnetic fields [2–4], this has important implications for the interpretation of the bound of eq. (2) for $(V-A)(V+A)$ interactions. In a recent paper [10] (see also [11,12]), we explored neutral (pseudo)scalar induced $M-\bar{M}$ oscillations. Since $(S \mp P)(S \pm P)$ and $(V \mp A)(V \pm A)$ operators are related by Fierz transform, we wish to study the magnetic field dependence of neutral scalar induced muonium conversion.

Exotic neutral (pseudo)scalar bosons $H$ and $A$ could have couplings [10],
\[ -\mathcal{L}_Y = \frac{f_H}{\sqrt{2}} \bar{\mu} e H + i \frac{f_A}{\sqrt{2}} \bar{\mu} \gamma_5 e A + h.c., \]  

(3)

while a discrete symmetry such as multiplicative \([1]\) electron number \(P_e\) \([13]\) forbids processes odd in number of electrons (plus positrons) like \(\mu \rightarrow e\gamma\) and \(\mu \rightarrow ee\bar{e}\). The resulting effective Hamiltonian responsible for \(M-\bar{M}\) conversion is,

\[ H_{M\bar{M}} = \frac{f_H^2}{2m_H^2} \bar{\mu} e \bar{e} - \frac{f_A^2}{2m_A^2} \bar{\mu} \gamma_5 e \bar{\mu} \gamma_5 e \]

\[ = \frac{1}{4} \left( \frac{f_H^2}{m_H^2} + \frac{f_A^2}{m_A^2} \right) (S^2 - P^2) + \frac{1}{4} \left( \frac{f_H^2}{m_H^2} - \frac{f_A^2}{m_A^2} \right) (S^2 + P^2), \]  

(4)

where \(S = \bar{\mu} e\) and \(P = \bar{\mu} \gamma_5 e\). Note that \(S^2 - P^2\) and \(S^2 + P^2\) contain \((S \mp P)(S \pm P)\) and \((S \mp P)(S \mp P)\) interactions, respectively. In the “\(U(1)\) limit” of \(f_H = f_A\) and \(m_H = m_A\) the subleading \((S \mp P)(S \mp P)\) terms are completely absent.

The matrix elements of eq. (4) and the accompanying phenomenology of eq. (3) have been discussed in ref. \([10]\). To explore magnetic field dependence of muonium conversion probabilities, consider the Hamiltonian for 1S muonium,

\[ \mathcal{H} = \mathcal{H}_0 + a s_e \cdot s_\mu - \mu_e \cdot B - \mu_\mu \cdot B + \mathcal{H}_{M\bar{M}}, \]  

(5)

where \(\mathcal{H}_0\) gives the 1S energy \(E_0 = -\alpha^2 m/2\), with reduced mass \(1/m = 1/m_e + 1/m_\mu\), \(a \approx 1.846 \times 10^{-5}\) eV is the 1S muonium hyperfine splitting, and \(\mu_e = -g_e \mu_B s_e\), \(\mu_\mu = g_\mu \mu_B m_\mu s_\mu\), where \(g_e \approx g_\mu \approx 2\) and \(\mu_B = e/(2m_e) \approx 5.788 \times 10^{-9}\) eV/Gauss is the Bohr magneton. Introducing the dimensionless parameters,

\[ X, Y \equiv \frac{\mu_B B}{a} \left( g_e \pm \frac{m_e}{m_\mu} g_\mu \right), \]  

(6)

we see that \(|Y|\) is just 1% smaller than \(|X|\). Ignoring \(\mathcal{H}_{M\bar{M}}\) for the moment, the four muonium Breit-Rabi energy levels \([14]\) are,

\[ E_M(1, \pm 1) = E_0 + \frac{a}{2} \left( \frac{1}{2} \pm Y \right), \]

\[ E_M(1, 0) = E_0 + \frac{a}{2} \left( -\frac{1}{2} + \sqrt{1 + X^2} \right), \]

\[ E_M(0, 0) = E_0 + \frac{a}{2} \left( -\frac{1}{2} - \sqrt{1 + X^2} \right), \]  

(7)
which correspond to the eigenstates

\[ |M; 1, \pm 1\rangle = |M; \uparrow\uparrow\rangle, |M; \downarrow\downarrow\rangle, \]
\[ |M; 1, \phantom{0} 0\rangle = c |M; \uparrow\downarrow\rangle + s |M; \downarrow\uparrow\rangle, \]
\[ |M; 0, \phantom{0} 0\rangle = -s |M; \uparrow\downarrow\rangle + c |M; \downarrow\uparrow\rangle, \quad (8) \]

where the magnetic field dependent “rotation” is

\[ s = \frac{1}{\sqrt{2}} \left[ 1 - \frac{X}{\sqrt{1 + X^2}} \right]^{\frac{1}{2}}, \quad c = \frac{1}{\sqrt{2}} \left[ 1 + \frac{X}{\sqrt{1 + X^2}} \right]^{\frac{1}{2}}. \quad (9) \]

We have labeled the Breit-Rabi energy levels with the weak field basis \(|M; F, m_F\rangle\), i.e. the corresponding zero \(B\) field (\(X, Y \rightarrow 0\) and \(s, c \rightarrow 1/\sqrt{2}\)) hyperfine states. The “uncoupled” basis of \(|M; m_{se}, m_{su}\rangle\) corresponds to the strong field limit of \(X, Y \gg 1\) and \(s \rightarrow 0\). In this limit, \(|M; 1, 0\rangle \rightarrow |M; \uparrow\downarrow\rangle\) and \(|M; 0, 0\rangle \rightarrow |M; \downarrow\uparrow\rangle\), and the electron and muon hyperfine spin-spin coupling is overwhelmed by the Zeeman effect.

For antimuonium, again ignoring the effect of \(H_{\tilde{M}\bar{M}}\), retaining the spin labels and with uncoupled basis \(|\tilde{M}; m_{se}, m_{su}\rangle\), one simply flips \(X \rightarrow -X, Y \rightarrow -Y\) and hence interchange \(s \leftrightarrow c\) in eqs. (8) and (9). The notable changes are

\[ E_{\bar{M}}(1, \mp 1) = E_M(1, \pm 1), \quad (10) \]

since for given spin, the antiparticle magnetic moments have flipped sign, and,

\[ |\tilde{M}; 1, 0\rangle = s |\tilde{M}; \uparrow\downarrow\rangle + c |\tilde{M}; \downarrow\uparrow\rangle, \]
\[ |\tilde{M}; 0, 0\rangle = -c |\tilde{M}; \uparrow\downarrow\rangle + s |\tilde{M}; \downarrow\uparrow\rangle. \quad (11) \]

We have the following energy differences between \(M\) and \(\tilde{M}\) eigenstates,

\[ E_M(1, \pm 1) - E_{\bar{M}}(1, \pm 1) = \pm a Y, \]
\[ E_M(1, \phantom{0} 0) - E_{\bar{M}}(1, \phantom{0} 0) = E_M(0, \phantom{0} 0) - E_{\bar{M}}(0, \phantom{0} 0) = 0, \]
\[ E_M(1, \phantom{0} 0) - E_{\bar{M}}(0, \phantom{0} 0) = -(E_M(0, \phantom{0} 0) - E_{\bar{M}}(1, \phantom{0} 0)) = a \sqrt{1 + X^2}. \quad (12) \]

The effect of \(H_{M\bar{M}}\) can be treated as a perturbation. Define generically
\[ \langle \tilde{M}|\mathcal{H}_{M\tilde{M}}|M\rangle = \frac{\delta}{2} \quad (13) \]

(for simplicity, we take \( \delta \) to be real) between any two Breit-Rabi \( M \) and \( \tilde{M} \) energy eigenstates, it was shown by Feinberg and Weinberg [1] that the time integrated probability for an initial muonium state to decay as antimuonium is

\[ P(\tilde{M}) = \frac{\delta^2}{2(\delta^2 + \Delta^2 + \lambda^2)}, \quad (14) \]

where \( \Delta = E_M - E_{\tilde{M}} \) is the energy difference, and \( \lambda \approx 2.996 \times 10^{-10} \text{ eV} \) is the muon decay rate. The physics is clear: muonium oscillation has to compete with muon decay and the damping from oscillations between two states that are very disparate in energy. The total transition probability is

\[ P_T(\tilde{M}) = \sum_{F,m_F} |c_{F,m_F}|^2 P^{(F,m_F)}(\tilde{M}) \quad (15) \]

where \( |c_{F,m_F}|^2 \) are the populations in muonium states of eq. (8), and \( P^{(F,m_F)}(\tilde{M}) \) are the probabilities for an initial \((F,m_F)\) muonium state to decay as antimuonium.

In principle \( P^{(1,0)}(\tilde{M}) \) and \( P^{(0,0)}(\tilde{M}) \) also contain the probabilities of initial \((1,0)\) or \((0,0)\) muonium states to decay as \((0,0)\) or \((1,0)\) antimuonium, respectively, since \((1,0)\) and \((0,0)\) are mixtures of unpolarized states. To show that these are vanishingly small, it is useful to notice the hierarchy

\[ \delta \ll \lambda \ll a, \quad (16) \]

the first of which follows from eq. (2) for any \( \mathcal{H}_{M\tilde{M}} \) model. Combining eqs. (12) and (14) we see that \((1,0) \rightarrow (0,0)\) and \((0,0) \rightarrow (1,0)\) transitions are extremely suppressed by hyperfine splitting even for \( B = 0 \), and need not be considered. Similarly, although \((1,\pm 1) \rightarrow (1,\pm 1)\) transitions contribute in \( B = 0 \) limit, they become rapidly suppressed even for rather weak magnetic fields [1], because of the mismatch between Zeeman energy levels for \( M \) vs. \( \tilde{M} \) polarized states with same \( m_F \).

In this work we will discuss only the relative magnetic field dependence of \( M-\tilde{M} \) conversion probabilities. Differences in coupling strength for various effective operators at zero \( B \)
field can be found in refs. [5–7,10]. Hence, we normalize effective interactions to the conversion probability due to eq. (1) at zero magnetic field, and in particular for equally populated (1/4 each) Breit-Rabi states (the latter condition would be removed at the end). Thus, the zero field total transition probability is

\[ P_T(\bar{M}; B = 0) \approx \frac{\bar{\delta}^2}{2\lambda^2} = 2.56 \times 10^{-5} \left( \frac{G_{M \bar{M}}}{G_F} \right)^2, \tag{17} \]

which defines the parameter \( \bar{\delta} \).

For the \((V - A)(V - A)\) case one basically multiplies the r.h.s. (right hand side) of eq. (13) by \( \delta_{m_e m_e} \delta_{m_e m_\mu} \), hence

\[
\langle \bar{M}; 1, \pm 1 | \mathcal{H}_{MM} | M; 1, \pm 1 \rangle = \frac{\bar{\delta}}{2},
\]

\[
\langle \bar{M}; 1, 0 | \mathcal{H}_{MM} | M; 1, 0 \rangle = \langle \bar{M}; 0, 0 | \mathcal{H}_{MM} | M; 0, 0 \rangle = \frac{\bar{\delta}}{2\sqrt{1 + X^2}}, \tag{18}
\]

and all other matrix elements vanish. One therefore finds the result [15]

\[
P^{(1, \pm 1)}(\bar{M}) = \frac{\bar{\delta}^2}{2(\bar{\delta}^2 + a^2 Y^2 + \lambda^2)},
\]

\[
P^{(1, 0)}(\bar{M}) = P^{(0, 0)}(\bar{M}) = \frac{\bar{\delta}^2}{2(\bar{\delta}^2 + \lambda^2 (1 + X^2))}, \tag{19}
\]

The magnetic field dependence for \( P_T(\bar{M}) \), as well as the separate probabilities \( P^{(1, \pm 1)}(\bar{M}) \), \( P^{(1, 0)}(\bar{M}) \) and \( P^{(0, 0)}(\bar{M}) \) of eq. (19), are plotted as “+” symbols in Figs. 1–4, respectively. The behavior is readily understood. Because of the \( aY \) energy splitting, the suppression of \((1, \pm 1)\) modes sets in with \( B \) field of just a few cG, and they become quenched for 0.1 G or higher. The \( m_F = 0 \) “unpolarized” modes are oblivious to the magnetic field until \( X \) becomes appreciable, i.e. for \( B \sim a/2\mu_B \sim 1kG \), and get quenched by fields of 1 Tesla or higher. The scale difference for Fig. 4 would be discussed shortly. The suppression in \( P(\bar{M}) \) has been taken into account in the experimental limit of eq. (2) for the interaction of eq. (1).

We have checked and confirmed the result for the \((V - A)(V + A)\) case [9],

\[
P^{(1, \pm 1)}(\bar{M}) = \frac{\bar{\delta}^2}{6(\bar{\delta}^2 + a^2 Y^2 + \lambda^2)},
\]
\[ P^{(1, \ 0)}(\bar{M}), \ P^{(0, \ 0)}(\bar{M}) = \frac{(2 \mp \frac{1}{\sqrt{1+X^2}})^2 \delta^2}{6 \left[ (2 \mp \frac{1}{\sqrt{1+X^2}})^2 \delta^2 + \lambda^2 \right]}. \]  

The results are also plotted in Figs. 1–4 as “×” symbols. Conversion in (0, 0) mode is the most prominent, but gets suppressed by up to a factor of 4/9 when the magnetic field goes beyond \( \sim \) 1kG. The (1, ±1) modes are quenched by magnetic fields of 0.1G or higher, just like in the \((V - A)(V - A)\) case, but \(P^{(1,0)}(\bar{M})\) actually grows with B field around 1kG, and partially compensates for the drop in \(P^{(0,0)}\).

We now state the results for the (pseudo)scalar induced interaction case. Details would be given elsewhere. For purely scalar interactions \((f_A = 0)\), we find

\[ P^{(1, \pm 1)}(\bar{M}) = \frac{\delta^2}{2(\delta^2 + a^2Y^2 + \lambda^2)}, \]
\[ P^{(1, \ 0)}(\bar{M}) = P^{(0, \ 0)}(\bar{M}) = \frac{\delta^2}{2(\delta^2 + \lambda^2)}. \]  

(21)

Thus, aside from the familiar quenching of the \((1, \pm 1)\) states, the \((1, 0)\) and \((0, 0)\) states are completely insensitive to magnetic fields. For purely pseudoscalar interactions \((f_H = 0)\), we find

\[ P^{(1, \pm 1)}(\bar{M}) = 0, \]
\[ P^{(1, \ 0)}(\bar{M}), \ P^{(0, \ 0)}(\bar{M}) = \frac{(\mp 1 + \frac{1}{\sqrt{1+X^2}})^2 \delta^2}{2 \left[ (\mp 1 + \frac{1}{\sqrt{1+X^2}})^2 \delta^2 + \lambda^2 \right]}. \]  

(22)

In this case, muonium conversion occurs solely in the \((0, 0)\) mode for zero magnetic field. Conversion in the \((1, 0)\) mode starts to grow from zero for field strengths beyond \( \sim \) 1kG, and partially compensates for the drop, by a factor of 4, in transition probability in the \((0, 0)\) mode. We plot the results again in Figs. 1–4, with solid and dashed lines representing scalar and pseudoscalar case, respectively.

The results for \((S - P)(S - P)\) and \((S - P)(S + P)\) operators can be similarly obtained. With the same normalization conditions as described above, the results are also given in Figs. 1–4 for sake of comparison, with open circles representing the \((S - P)(S + P)\) case and open boxes representing the \((S - P)(S - P)\) case. Note that according to eq. (4), the
\( (S - P)(S - P) \) operators should be subdominant compared to \( (S - P)(S + P) \) operators. Pure \( (S \mp P)(S \pm P) \) operators corresponds to complex neutral scalars \([10]\), where the sneutrino \( \tilde{\nu}_\tau \) in SUSY models with \( R \)-parity breaking \([17]\) as a special case. It is evident from Figs. 1–4 that the combinations \( (S - P)(S - P) \) and \( (S - P)(S + P) \) behave in the same way as \( (V - A)(V - A) \) and \( (V - A)(V + A) \), respectively. In the latter case, the two operators are related to each other by a Fierz transform. For the former case, although the operators can not be related to each other by a Fierz transform, the matrix elements are always in same proportion, which comes as a consequence of the nonrelativistic limit.

Turning to discussions, we note that the assumption of equally populated Breit-Rabi levels is not a valid one, since this is determined by the muonium formation process and the magnetic field strength. However, as a consequence of this assumption, the results in eqs. (19-22) and hence Figs. 2–4 all have an artificial factor of 4. To illustrate the effect of differently populated Breit-Rabi levels, normalizing again to the \( (V - A)(V - A) \) case at zero \( B \) field, we take muonium states to be populated as \([4]\) 32\%, 35\%, 18\% and 15\%, respectively, for \( (F, m_F) = (0, 0), (1, +1), (1, 0), (1, -1) \), and plot the results for \( P_T(\bar{M}) \) (i.e. analogous to Fig. 1) in Fig. 5. The \( (V - A)(V - A) \) and \( (S - P)(S - P) \) results are unchanged, since the conversion matrix elements are the same for all modes. The purely scalar \( (SS) \) case is also unchanged, since \( |c_{10}|^2 + |c_{00}|^2 = |c_{1+1}|^2 + |c_{1-1}|^2 = 50\% \) is the same as the equally populated case. The \( (V - A)(V + A) \), \( (S - P)(S + P) \) and \( PP \) cases are somewhat modified from the equally populated case, but the difference for 1kG field is rather slight.

Let us summarize our findings. Purely scalar \( (SS) \) induced \( M-\bar{M} \) transitions in polarized (i.e. \( (F,m_F) = (1, \pm 1) \)) modes are quenched by magnetic fields of 0.1 Gauss or higher, but conversion in unpolarized states (i.e. \( (1,0) \) and \( (0,0) \)) are independent of magnetic field strength. Purely pseudoscalar \( (PP) \) interactions do not induce conversion in \( (1, \pm 1) \) states, but exhibit compensating effects in the \( (1,0) \) and \( (0,0) \) channels, similar to the \( (V - A)(V + A) \) case. Interactions of \( (S \mp P)(S \pm P) \) and \( (V \mp A)(V \pm A) \) form have the same magnetic field dependence since they are related by Fierz transform, while \( (S \mp P)(S \mp P) \)
and \((V ± A)(V ± A)\) interactions have the same field dependence because of proportional conversion matrix elements.

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FIGURES

FIG. 1. Magnetic field dependence of total muonium conversion probability $P_T(\bar{M})$ assuming $|c_{1,1}|^2 = |c_{1,0}|^2 = |c_{1,-1}|^2 = |c_{0,0}|^2 = 1/4$, and normalized to conversion strength of $(V - A)(V + A)$ interaction at zero magnetic field. Solid and dashed lines stand for $SS$ and $PP$ operators, respectively, while $\circ$, $\square$, $+$ and $\times$ stand for $(S - P)(S + P)$, $(S - P)(S - P)$, $(V - A)(V + A)$ and $(V - A)(V - A)$ cases, respectively.

FIG. 2. $P^{(1,\pm 1)}(\bar{M})$ vs. magnetic field with same assumption as Fig. 1.

FIG. 3. $P^{(1,0)}(\bar{M})$ vs. magnetic field with same assumption as Fig. 1.

FIG. 4. $P^{(0,0)}(\bar{M})$ vs. magnetic field with same assumption as Fig. 1. The same as Fig. 1 except $|c_{1,1}|^2 = 0.35$, $|c_{1,0}|^2 = 0.18$, $|c_{1,-1}|^2 = 0.15$ and $|c_{00}|^2 = 0.32$. 
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