$\mu^+ e^- \leftrightarrow \mu^- e^+$ Transitions via Neutral Scalar Bosons

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With $\mu \rightarrow e\gamma$ decay forbidden by multiplicative lepton number conservation, we study muonium–antimuonium transitions induced by neutral scalar bosons. Pseudoscalars do not induce conversion for triplet muonium, while for singlet muonium, pseudoscalar and scalar contributions add constructively. This is in contrast to the usual case of doubly charged scalar exchange, where the conversion rate is the same for both singlet and triplet muonium. Complementary to muonium conversion studies, high energy $\mu^+ e^- \rightarrow \mu^- e^+$ and $e^- e^- \rightarrow \mu^- \mu^-$ collisions could reveal spectacular resonance peaks for the cases of neutral and doubly charged scalars, respectively.

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The interest in muonium–antimuonium ($M-\bar{M}$) conversion dates back to a suggestion by Pontecorvo \[1\], which pointed out the similarity between the $M-\bar{M}$ and $K^0-\bar{K}^0$ systems. Feinberg and Weinberg \[2\] noted further that $M-\bar{M}$ conversion is allowed by conservation of multiplicative muon number — muon parity — but forbidden by the more traditional additive muon number. It thus provides a sensitive test of the underlying conservation law for lepton number(s) and probes physics beyond the standard model. One advantage of studying $M-\bar{M}$ conversion is that, once the effective four-fermion Hamiltonian is given, everything is readily calculable since it involves just atomic physics. The experiment is quite challenging, however, while on the theoretical front, it has attracted less attention from model builders compared to decay modes like $\mu \rightarrow e\gamma$ which are in fact forbidden by the multiplicative law.

The effective Hamiltonian is traditionally taken to be of $(V-A)(V-A)$ form

$$H_{M\bar{M}} = \frac{G_{M\bar{M}}}{\sqrt{2}} \bar{\mu}\gamma_{\lambda}(1-\gamma_5)e\bar{\mu}\gamma_{\lambda}(1-\gamma_5)e + \text{H.c.},$$  \hspace{1cm} (1)$$

and experimental results are given \[3\] as upper limits on $R_g \equiv G_{M\bar{M}}/G_F$, where $G_F$ is the Fermi constant. The present limit is $R_g < 0.16$ \[4\]. This would soon be improved to $10^{-2}$ level \[5\] by an ongoing experiment \[6\] at PSI, with the ultimate goal of $10^{-3}$.

Explicit models that lead to effective interactions of eq. (1) were slow to come by. In 1982, Halprin \[7,8\] pointed out that in left-right symmetric (LRS) models with Higgs triplets, doubly charged scalars $\Delta^{--}$ can mediate $M-\bar{M}$ transitions at tree level in the $t$-channel (Fig. 1(a)). The effective interaction, after Fierz rearrangement, can be put in $(V \pm A)(V \pm A)$ form of eq. (1). This not only encouraged experimental interests \[3\], it also stimulated theoretical work \[9\]. In particular, Chang and Keung \[10\] give the conditions for a generic model. These work together gives one the impression that doubly charged scalar bosons may be the only credible source for inducing $M-\bar{M}$ transitions. However, in a recent model \[11\] for radiatively generating lepton masses from multiple Higgs doublets, it was pointed out in passing that the flavor-changing neutral Higgs bosons responsible for mass generation could also mediate $M-\bar{M}$ conversion. A remnant $Z_2$ symmetry serves the function
analogous [10] to Feinberg-Weinberg’s muon parity that forbids $\mu \rightarrow e\gamma$ transitions, while the effective four-fermion operators responsible for $M-\bar{M}$ transition are not of the form of eq. (1). In this paper we explore neutral scalar induced $M-\bar{M}$ oscillations [12] in the general case. Constraints from $g-2$ and $e^+e^- \rightarrow \mu^+\mu^-$ scattering data are studied. We point out that, complementary to muonium studies, high energy $\mu^+e^- \rightarrow \mu^-e^+$ and $e^-e^- \rightarrow \mu^-\mu^-$ collisions could clearly distinguish between (flavor changing) neutral and doubly charged scalar bosons.

Consider neutral scalar and pseudoscalar bosons $H$ and $A$, with the interaction,

$$-\mathcal{L}_Y = \frac{f_H}{\sqrt{2}} \bar{\mu}e H + i \frac{f_A}{\sqrt{2}} \bar{\mu}\gamma_5 e A + H.c.$$  \hspace{1cm} (2)

Imposing a discrete symmetry $P_e$ [10] such that the electron as well as $H, A$ fields are odd while the muon field is even, processes odd in number of electrons (plus positrons) like $\mu \rightarrow e\gamma$ and $\mu \rightarrow ee\bar{e}$ are forbidden. Namely, scalar bosons may not possess flavor diagonal and nondiagonal couplings at the same time. $P_e$ is nothing but a variation of the multiplicative muon number of Feinberg and Weinberg [4]. The interaction of eq. (2) induces (Figs. 1(b) and 1(c)) the effective Hamiltonian

$$\mathcal{H}_{S,P} = \frac{f_H^2}{2m_H^2} \bar{\mu}e \bar{\mu}e - \frac{f_A^2}{2m_A^2} \bar{\mu}\gamma_5 e \bar{\mu}\gamma_5 e,$$ \hspace{1cm} (3)

at low energy which mediates $M-\bar{M}$ conversion. The conversion matrix elements for $S^2$ and $P^2$ operators ($S, P$ stand for $\bar{\mu}e$ and $\bar{\mu}\gamma_5 e$ densities) are

$$\langle \bar{M}(F=0)|S^2|M(F=0)\rangle = +\frac{2}{\pi a^3}, \quad \langle \bar{M}(F=1)|S^2|M(F=1)\rangle = -\frac{2}{\pi a^3},$$ \hspace{1cm} (4)

$$\langle \bar{M}(F=0)|P^2|M(F=0)\rangle = -\frac{4}{\pi a^3}, \quad \langle \bar{M}(F=1)|P^2|M(F=1)\rangle = 0,$$ \hspace{1cm} (5)

where $F$ is the muonium total angular momentum, while $a$ is its Bohr radius. Thus, only scalars induce muonium conversion in the spin triplet state, while for singlet muonium, the effect of scalar and pseudoscalar channels add constructively. Note that for $(V \pm A)^2$ interactions of eq. (1), we always get $8G_{M\bar{M}}/\pi a^3$ for both singlet and triplet muonium [2]. One clearly sees that separate measurements of singlet vs. triplet $M-\bar{M}$ conversion

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probabilities can distinguish between neutral scalar, pseudoscalar and doubly charged Higgs induced interactions.

In practice, $M$ is formed as a mixture of triplet and singlet states. It is crucial whether the (anti)muon decays in the presence of magnetic fields. Any sizable field strength lifts the degeneracy of $M$–$ar{M}$ for $F = 1$, $m_F = \pm 1$ states, hence effectively “quenches” $M$–$ar{M}$ conversion. This is normally the case under realistic conditions, but experiments correct for this and report $G_{M\bar{M}}$ (or $R_g$) for zero $B$ field. It is important to note, however, that in so doing, one inadvertently ignores the possible differences in the neutral (pseudo)scalar case. Let us take the example of the ongoing PSI experiment $[6]$. Muonium is formed and stays in the presence of 1kG magnetic field. In this case, muonium states are populated as 32%, 35%, 18% and 15%, respectively, for $(F, m_F) = (0, 0), (1, +1), (1, 0), (1, -1)$. Only the $m_F = 0$ modes are active for muonium conversion, hence the effective triplet probability comes only from $|c_{1,0}|^2 = 18\%$, down from 68%. For $(V \pm A)^2$ interactions, one simply corrects for a factor of 1/2 reduction. For our case of neutral scalar induced interactions, the experimental limit on $G_{M\bar{M}}$ relates to scalar couplings as

$$
G_{M\bar{M}}^{\text{expt.}} = \frac{1}{8} \sqrt{\frac{|c_{0,0}|^2 \left( \frac{f_H^2}{m_H^2} + 2 \frac{f_A^2}{m_A^2} \right)^2 + |c_{1,0}|^2 \left( -\frac{f_H^2}{m_H^2} \right)^2}.
$$

(6)

Several cases are of interest: (a) $f_A = 0$; (b) the "U(1) limit" of $m_A = m_H$ ($H$ and $A$ form a complex neutral scalar), with $f_A = f_H$; (c) $f_H = 0$ (pseudoscalar only). For case (a), the result is rather similar to eq. (1). For case (b), constructive interference strongly enhances the effect in singlet channel. For case (c), only the singlet $(0, 0)$ part is active. If the PSI experiment will attain $[5]$ the limit of $R_g < 10^{-2}$ without observing $M$–$ar{M}$ conversion, eq. (6) would imply the bounds

$$
f^2/m^2 \lesssim (0.9, 0.4, 0.6) \times 10^{-6} \text{ GeV}^{-2},
$$

(7)

respectively, for the three cases, where $f/m$ stand for $f_H/m_H$ except for case (c).

Some other constraints on $\mathcal{H}_{S,P}$, such as the anomalous magnetic moments of the electron and muon, should be considered. Defining $a \equiv (g - 2)/2$, we find that
\[ \delta a_e \simeq -\frac{f^2}{16\pi^2} m_e \left( \frac{m_e}{3m^2} \mp \frac{3m_\mu}{2m^2} \mp \frac{m_\mu}{m^2} \ln \frac{m^2}{m^2} \right), \tag{8} \]

where \( \mp \) is for \( H \) or \( A \) contribution, respectively, while for \( a_\mu \) one interchanges \( e \longleftrightarrow \mu \).

Comparing experimental measurements \( \text{[3]} \) with QED prediction, we find \( \delta a_e^{\text{expt}} = (146 \pm 46) \times 10^{-12} \) and \( \delta a_\mu^{\text{expt}} = (27 \pm 69) \times 10^{-10} \). The effective bound from \( \delta a_e^{\text{expt}} \) on \( f^2/m^2 \) is of order \( G_F \), except for the \( U(1) \) limit case. In the latter case, cancellations between \( H \) and \( A \) lead to a much weaker limit. However, for muon \( g-2 \) the leading term (proportional to \( m_\mu^2 \)) comes from the first term of eq. (8) which does not suffer from \( H-A \) cancellation. Hence, it gives a bound of order \( 10 \, G_F \) for all cases. In any rate, these limits are considerably weaker than eq. (7).

An interesting constraint comes from high energy \( e^+e^- \rightarrow \mu^+\mu^- \) scattering cross sections, which probe the interference effects between the contact terms of eq. (2) (Fig. 1(b) in \( t \)-channel) and standard diagrams. For case (b), the effective contact interaction can be put in standard form \( \text{[13]} \) for compositeness search,

\[ \mathcal{H}_{ee\mu\mu} = \frac{f^2}{2m^2} (\bar{\mu}e \bar{\mu}e - \bar{\mu}\gamma_5 e \bar{\mu}\gamma_5 e) \]
\[ = \frac{g^2}{2\Lambda^2} \{ \bar{e}\gamma_\alpha R e \bar{\mu}\gamma^\alpha L \mu + \bar{e}\gamma_\alpha L e \bar{\mu}\gamma^\alpha R \mu \}, \tag{9} \]

where \( \Lambda \equiv \Lambda_{LR}^+ \). Setting \( g^2/(4\pi) = 1 \), the combined limit gives \( \Lambda(ee\mu\mu) > 2.6 \text{ TeV} \) \( \text{[13]} \), which translates to \( f^2/m^2 < 1.9 \times 10^{-6} \text{ GeV}^{-2} \). This can be converted to a limit on \( M-\bar{M} \) conversion by assuming eq. (6),

\[ G_{M\bar{M}} < 0.06 \, G_F, \tag{10} \]

which is better than present \( \text{[4]} \) \( M-\bar{M} \) conversion bound of \( R_g < 0.16 \), but somewhat weaker than the \( 10^{-2} \) bound expected soon at PSI \( \text{[3]} \).

In the model of ref. \( \text{[1]} \), scalar interactions of the type of eq. (2) were used to generate charged lepton masses iteratively order by order, via effective one loop diagrams with lepton seed masses from one generation higher. To be as general as possible, we are not concerned with the generation of \( m_\mu \) from \( m_\tau \) here. However, in analogy to the softly broken \( Z_8 \)
symmetry of ref. [11], some discrete symmetry can be invoked to forbid electron mass at
tree level but allow it to be generated by $m_\mu$ via one loop diagrams as shown in Fig. 2.
Since $m_{H,A} \gg m_\mu$, we have

$$\frac{m_e}{m_\mu} \approx \frac{f^2}{32\pi^2} \log \frac{m_H^2}{m_A^2}. \quad (11)$$

Note that $f_H = f_A = f$ is necessary for divergence cancellation, hence in the $U(1)$ limit
[11] of $m_A = m_H$ the mass generation mechanism is ineffective. We see that, because the
factor of $1/32\pi^2 \sim 1/300$ is already of order $m_e/m_\mu$, if $m_A \neq m_H$ but are of similar order
of magnitude, in general we would have $f \sim 1$. This looks attractive for scalar masses far
above the weak scale since one could have large Yukawa couplings but at the same time
evade the bound of eq. (7). However, in the more ambitious model of ref. [11], radiative
mass generation mechanism is pinned to the weak scale, namely, Higgs boson masses cannot
be far above TeV scale for sake of naturalness. In this case, although eq. (11) still looks
attractive and is a simplified version of the more detailed results of ref. [11], with $f \sim 1$ and
$m_H, m_A \lesssim$ TeV, the bound of eq. (7) cannot be satisfied. We thus conclude that the bound
of eq. (7), expected soon from PSI, will rule out the possibility of radiatively generating $m_e$
solely from $m_\mu$ via one loop diagrams, if the lepton number changing neutral scalar bosons
are of weak scale mass. A model where $m_e$ dominantly comes from $m_\tau$ at one loop level,
with a minor contribution from $m_\mu$, would be presented elsewhere [14].

If $M-M$ conversion is observed, one would certainly have to make separate measurements
in singlet vs. triplet states to distinguish between the possible sources. Complementary
to this, one could explore signals at high energies. It was pointed out already by Glashow
[15] the connection between the studies of $e^-e^- \to \mu^-\mu^-$ collisions and $M-M$ conversion.
Indeed, shortly after the first $M-M$ experiment [16], studies of $e^-e^-$ collisions at SLAC
improved the limit on $G_{M\bar{M}}$ by a factor of 10 [17]. Although such efforts have not been
repeated, it has been stressed recently by Frampton [18] in the context of dilepton gauge
bosons [19]. It is clear that if $\Delta^{-}$ exists, it would appear as a resonance peak in energetic
$e^-e^- \to \mu^-\mu^-$ collisions.
In contrast, it has rarely been mentioned that $\mu^+ e^- \rightarrow \mu^- e^+$ collisions may also be of great interest. Even for $\Delta^{--}$ bosons, the cross section can be sizable for $\sqrt{s} \sim m_\Delta$. However, if neutral scalars that mediate $M-\bar{M}$ conversion exist, and the masses are of order TeV or below, one would have spectacular s-channel resonances in $\mu^\pm e^\mp$ collisions! Even the non-observation of $M-\bar{M}$ conversion does not preclude this possibility. Let us assume that PSI would not observe $M-\bar{M}$ conversion at $10^{-2}$ level, hence $f^2/m^2$ is bound by eq. (7). Assuming just a single scalar boson $H$ (case (a)) that saturates such a bound, and that $H \rightarrow \mu^\pm e^\mp$ only, we plot in Fig. 3 the cross section $\sigma(\mu^+ e^- \rightarrow \mu^- e^+)$ vs. $\sqrt{s}$ for $m_H = 0.25, 0.5, 1$ and $2$ TeV. The result for $\Delta^{--}$ constrained by $G_{MM} \lesssim 10^{-2}$ is also shown in Fig. 3 as dashed lines for similar masses. Note that for $f = 0.1 - 2$, which is the plausible range for Yukawa couplings advocated in ref. [11,14], eq. (7) implies that the lower bound for $m_H$ ranges between 100 GeV – 2 TeV. For $e^- e^- \rightarrow \mu^- \mu^-$ collisions, the curves are rather similar, with the role of $H$ and $\Delta^{--}$ interchanged. It is clear that $\mu^+ e^-$ or $e^- e^-$ colliders in the few hundred GeV to TeV range have the potential of observing huge cross sections, and could clearly distinguish between $H$ and $\Delta^{--}$.

The development of $\mu^+ \mu^-$ colliders have received some attention recently [20]. Perhaps one could also consider the $\mu^\pm e^\mp$ collider option, especially if one could utilize existing facilities. As muons are collected via $\pi \rightarrow \mu$ decay, existing accelerator complexes that have both electron and proton facilities, such as CERN or HERA, are preferred. Since $\mu^+$ is easier to collect and cool, while $e^-$ requires no special effort, $\mu^+ e^-$ collisions should be easier to perform. For example, take $E_e$ to be the LEP II beam energy of 90 GeV, if intense 200 GeV – 7 TeV $\mu^+$ beams could be produced, one could attain $\sqrt{s} \simeq 190$ GeV – 1.1 TeV. Compared with problems like $\mu$ decay before collision for $\mu^+\mu^-$ colliders [20], $\mu^+ e^-$ events in $\mu^+ e^-$ collisions have practically no background. Future linear colliders should be able to span an even wider energy range, perhaps performing $e^- e^-$, $\mu^\pm e^\mp$, $\mu^+ \mu^-$ as well as $e^- e^+$ collisions.

Some discussion is in order. Neutral scalars with flavor changing couplings may appear to be exotic [21]. However, with multiplicative lepton number, one evades the bounds from
\( \mu \to e\gamma \) decay and the like. In this light, we note that any model with more than one Higgs doublet in general would give rise to flavor changing neutral scalars. Second, the couplings of eq. (2) demand that \( H \) and \( A \) carry weak isospin, hence they must have charged partners. These charged scalars can induce the so-called “wrong neutrino” decay \( \mu^- \to e^-\nu_e\bar{\nu}_\mu \). Third, the conversion matrix elements for \((S\pm P)^2\) part of eq. (3) can be related to eq. (1), but the \((S\pm P)(S\mp P)\) parts are related to \((V\pm A)(V\mp A)\) operators, which were considered by Fujii et al. \[22\] in the context of dilepton gauge bosons. In general, \( M-\bar{M} \) conversion may have four different kind of sources: doubly charged scalar or vector bosons in \( t \)-channel, or neutral scalar or vector bosons in \( s \)- or \( t \)-channel. Dilepton gauge boson models are therefore of the second type. Neutral vector bosons would come from horizontal gauge symmetries, but models are somewhat difficult to construct \[23\]. Detailed measurements of singlet vs. triplet \( M-\bar{M} \) conversion, as well as high energy \( \mu^\pm e^\mp \to \mu^\mp e^\pm \) and \( e^-e^- \to \mu^+\mu^- \) collisions should be able to identify the actual agent for these lepton number violating interactions. Fourth, in supersymmetric theories containing \( R \)-parity violating terms \[24\], \( s \)-channel \( \tilde{\nu}_\tau \) (\( \tau \) sneutrinos, a kind of neutral scalar) exchange could also induce \( M-\bar{M} \) conversion, resulting in \((S - P)(S + P)\) operators.

Let us summarize the novel features of this work. We point out that neutral (pseudo)scalars may well induce muonium–antimuonium transitions, something that has been largely neglected in the literature. All one needs is to invoke multiplicative lepton number rather than adhering to the traditional but more restrictive additive lepton number conservation. In this way, stringent limits from \( \mu \to e\gamma \) decay, etc., are evaded. The induced operators differ from the usual \((V - A)(V - A)\) form, and care has to be taken when one interprets experimental limits. In particular, measuring \( M-\bar{M} \) conversion strength in both singlet and triplet muonia can distinguish between different interactions. A limit of \( G_{M\bar{M}} < 10^{-2} \ G_F \), expected soon at PSI, would rule out the possibility of radiatively generating \( m_e \) solely from \( m_\mu \) at one loop order via neutral scalar bosons with weak scale mass. Complementary to \( M\bar{M} \) studies, high energy \( \mu^+e^- \to \mu^-e^+ \) collisions may reveal resonance peaks for flavor changing neutral scalars, while the more widely known doubly charged scalar
would appear as resonances in $e^- e^- \rightarrow \mu^- \mu^-$ collisions.

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FIGURES

FIG. 1. Diagrams for $\mu^+e^- \rightarrow \mu^-e^+$ transitions via (a) doubly charged scalar $\Delta^{--}$; and (b), (c) neutral (pseudo)scalars $H, A$.

FIG. 2. One loop diagram for $m_e$ generation.

FIG. 3. $\sigma(\mu^+e^- \rightarrow \mu^-e^+)$ vs. $\sqrt{s}$ for $m_H = 0.25, 0.5, 1, 2$ TeV. Only $H \rightarrow \mu^\pm e^\mp$ is taken into account for $\Gamma_H$, with Yukawa couplings saturating $f^2_H/m^2_H \lesssim 0.9 \times 10^{-6}$ GeV$^{-2}$. Analogous bounds for the case of $\Delta^{--}$ is shown as dashed lines.
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