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Astrophysical Bound on the Majoron-Higgs Boson Coupling

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Abstract

We show that the coupling of the "standard" Higgs boson to Majorons, that could lead to a very fast decay of the neutral Higgs scalar to invisible modes, can be bounded using astrophysical arguments. We discuss the relevance of this bound for low-energy phenomenology related to majoron production. The bound so obtained may also jeopardize the stability of the VEV hierarchy in the doublet and triplet majoron models if the mass of the top quark is less than the W mass. A similar analysis applies to any model which exhibits Goldstone - or pseudo-Goldstone - bosons in the spectrum.

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The idea that a natural framework for a solution of the standing solar neutrino problem [?] through matter enhanced neutrino oscillations [?] may be provided by the spontaneous breaking of lepton number has been recently investigated by various authors [?,?]. In particular, a simple extension of the standard electroweak model in which the majoron [?,?], the Goldstone boson associated with the spontaneous braking of lepton number, belongs to an $SU(2)_L$ *doublet* has been proposed by the present authors [?] as a predictive model for solar neutrino oscillations. Due to its minimality the model is in fact sharply constrained by present phenomenology, thus allowing for a test of solar neutrino oscillation parameters in conventional experiments. Among them, the new scalar contributions to neutrino-like Z^0 decays and to $\mu \to e\gamma$ provide two of the most sensitive tests [?].

In more recent papers we emphasized the relevance of the majoron-Higgs boson coupling for some topical phenomenological issues, as a possibly large contribution to the $K \to \pi^+ + \textit{nothing}$ decay [?] or an elusive decay to majorons as the main decay mode for the "standard" Higgs boson over a large mass range [?]. In this respect it is particularly interesting that the study of the stability of the vacuum-expectation-value (VEV) hierarchy, present in general in non-singlet majoron models, leads, at the one-loop level to a relation among scalar couplings which limits, as a function of the top mass, the arbitrariness of the majoron-Higgs coupling [?].

In this paper we analize the consequences of the presence of the Higgs coupling to majorons on the evolution of helium burning red giant stars. It is known that the emission of light and weakly interacting particles in red giant stars may lead to a too rapid cooling of the core, thus preventing helium ignition [?]. This generally leads to strong limits on the rate of the processes associated with their production. In particular, for non singlet majoron models these considerations are at the origin of the strong bound on the lepton breaking

VEV v ($v \lesssim 10\ KeV$ [?]), related to majoron emission via the Compton process $\gamma + e \to e + J$ (the coupling of the majoron J to electrons is proportional to the mixing with the standard doublet and, therefore, to the lepton breaking VEV). The process we consider here is the analogous Compton emission in which *two* majorons are produced through the exchange of a virtual Higgs boson which couples to electrons via the standard Yukawa coupling. Due to the existing bound on v we may a priori expect this process to be highly competitive with the single majoron emission and therefore lead to a sharp bound on the Majoron-Higgs boson coupling. Our analysis shows indeed that the astrophysical constraint is sufficient to limit the Higgs boson mediated contribution to the rare process $K^+ \to \pi^+ +$ *nothing* (through the JJ or $\rho_L \rho_L$ modes), which could be by far the dominant one [?], to be of the same order of the gauge boson mediated contributions. The same bound, when applied to the question of the stability of the VEV hierarchy under radiative corrections in non-singlet majoron models, may require, for the perturbative consistency of the models, a top mass heavier than the W mass. Affected by the bound is also the size of the contribution to the width of the neutral Higgs boson related to "invisible" majoron modes. It is wortwhile to remark that the phenomenological issues related to the majoron-Higgs boson coupling are generally relevant for models which exhibit Goldstone or pseudo-Goldstone bosons, other than the majoron. The "invisibile" axion of Dine, Fischler and Srednicki [?] provides an example. The analysis and most of the conclusions of this paper, which explicitely refer to the cases of the triplet and doublet majoron models, are easily extended to other models.

The most general form of the Higgs potential in models with Goldstone bosons exhibits quartic terms that couple the standard Higgs doublet to the new multiplets introduced in the model. These terms are in general required by the renormalizability of the theory and after symmetry breaking lead to a coupling of the standard Higgs boson to a pair of Goldstone bosons proportional to the vacuum expectation value (VEV) of the standard doublet. In the

case of the doublet and triplet majoron models the relevant coupling can be written as [?]

$$\mathcal{L}_H \simeq -\frac{1}{\sqrt{2}}\lambda_3 u(J^2+\rho_L^2)H \tag{1}$$

where $u \simeq 174\ GeV$ is the standard VEV and ρ_L is the light scalar partner of the majoron whose mass is of the order or smaller than the lepton breaking VEV v ($m_{\rho_L} \lesssim 10\ KeV$). For the time being, we take the conservative approach of considering only majoron emission, since the Compton emission of a pair of ρ_L, otherwise analogous to JJ, may be kinematically disfavoured.

In order to evaluate the energy loss due to the process depicted in fig. 1, we have to evaluate the following expression

$$\epsilon = \frac{1}{\rho}\int dn_e \int dn_\gamma \frac{1}{4E_e E_\gamma}\int d\Phi_3 (1-f(E_e'))\sum_{pol}|\mathcal{M}|^2(E_1+E_2) \tag{2}$$

where ϵ represents the energy emission per unit mass at temperature T due to the production of two majorons of energy E_1 and E_2 respectively, $\mathcal{M}$ being the S-matrix element for the process considered. The electron and photon number densities dn_e and dn_γ are given respectively by

$$dn_e = d^3p\frac{2}{(2\pi)^3}\frac{1}{e^{(E-\mu)/T}+1} \tag{3}$$

$$dn_\gamma = d^3p\frac{2}{(2\pi)^3}\frac{1}{e^{E/T}-1} \tag{4}$$

where μ is the electron chemical potential and we have set the Boltzman constant to unit. It is worth noting that for the densities and temperatures typical of a red giant star we have $n_{e^+} \ll n_e$ [?] and therefore we may neglect the energy loss due to the Compton scattering over positrons. The three body phase space $d\Phi_3$ is given, following the notation of fig. 1, by

$$d\Phi_3 = (2\pi)^4\delta^4(p_\gamma+p_e-p_e'-q_1-q_2)\frac{d^3p_e'}{(2\pi)^3 2E_e'}\frac{d^3q_1}{(2\pi)^3 2E_1}\frac{d^3q_2}{(2\pi)^3 2E_2} \tag{5}$$

Finally, the factor $1 - f(E_e')$, where $f(E) \equiv (e^{(E-\mu)/T} + 1)^{-1}$, represents the Pauli blocking effect, which, for electrons in a degenerate regime may substantially reduce the efficiency of the reaction.

The expression in eq. (??) can be greatly simplified if we consider that for temperatures $T \simeq 10\ KeV$ and densities $\rho \sim 10^4\ g/cm^3$, typical of red giant stars, we may use, to a good approximation, the nonrelativistic and nondegenerate limit for the electron distribution. In this limit we may replace $E_1 + E_2$ by E_γ, neglect the Pauli blocking factor and write the expression for the emissivity directly in terms of the cross section for the process under consideration. After some trivial integrations one obtains

$$\epsilon = \frac{n_e}{\rho} \frac{1}{\pi^2} \int_0^\infty dE_\gamma \frac{E_\gamma^3}{e^{E_\gamma/T} - 1} \sigma(E_\gamma) \tag{6}$$

The electron number density n_e is related for a neutral plasma to the proton number density n_p by $n_e - n_{e^+} = n_p$. In our case $n_e \simeq n_p$, and considering that the chemical composition of the star is dominated by Hydrogen and Helium we can write the electron number density as $n_e = (\rho/2m_N)(1 + X_H)$, where X_H is the fractional mass abundance of Hydrogen and m_N is the nucleon mass.

The transition amplitude for the process $\gamma e \to eJJ$ is readily written from the diagrams in fig. 1 :

$$\mathcal{M} = \frac{2G_F e}{\sqrt{2}} \left(\frac{\lambda_3}{2\lambda_1}\right) m_e \bar{u}(p_e') \{ \frac{1}{\not{p}_e + \not{p}_\gamma - m_e} \gamma^\mu + \gamma^\mu \frac{1}{\not{p}_e' - \not{p}_\gamma - m_e} \} u(p_e) \epsilon_\mu(p_\gamma) \tag{7}$$

where we used the relations $m_H = 4\lambda_1 u^2$ and $G_F/\sqrt{2} = 1/4u^2$ and e is the conventionally defined proton charge. After averaging and summing over polarizations we obtain

$$\overline{\sum_{pol}} |\mathcal{M}|^2 = \left(\frac{G_F m_e e}{\sqrt{2}}\right)^2 \left(\frac{\lambda_3}{2\lambda_1}\right)^2 \left\{ 4 \left(\frac{m^2 - u}{s - m^2} + \frac{s - m^2}{m^2 - u} - 2 \right) \right.$$

$$\left. + 32 \left(\frac{m^2}{m^2 - u} - \frac{m^2}{s - m^2} \right) - 32 \left(\frac{m^2}{m^2 - u} - \frac{m^2}{s - m^2} \right)^2 \right.$$

$$+8s_2\left(\frac{1}{s-m^2}-\frac{1}{m^2-u}+\frac{m^2}{(m^2-u)^2}+\frac{m^2}{(s-m^2)^2}-\frac{6m^2-s_2}{(s-m^2)(m^2-u)}\right)\Bigg\} \tag{8}$$

where m is the electron mass and the invariants s, u and s_2 are defined as $s = (p_e + p_\gamma)^2$, $u = (p'_e - p_\gamma)^2$ and $s_2 = (q_1 + q_2)^2$. Integrating over the three body phase space we finally derive the (exact) expression

$$\sigma(z) = \frac{\alpha}{16(4\pi)^2}\left(\frac{\lambda_3}{2\lambda_1}G_F m_e\right)^2 \frac{f(z)}{z^2} \tag{9}$$

where $z \equiv (s - m_e^2)/m_e^2$ and the function $f(z)$ is defined as

$$f(z) = -\frac{190}{3} - \frac{80}{3}z - \frac{20}{3(1+z)} + \frac{2}{3(1+z)^2} - \frac{58}{9}z^2 + \left(56 + 16z + \frac{208}{3z} + \frac{8}{3}z^2\right)\log(1+z) \tag{10}$$

In the nonrelativistic limit $z \to 2E_\gamma/m_e \ll 1$ and $f(z) \to 8z^4/15$. Thus the required cross section in the nonrelativistic limit is given by

$$\sigma(E_\gamma) = \xi \frac{E_\gamma^2}{m_e^4} \tag{11}$$

where ξ is a dimensionless parameter proportional to the relevant couplings

$$\xi \equiv \frac{2\alpha}{15(4\pi)^2}\left(\frac{\lambda_3}{2\lambda_1}G_F m_e^2\right)^2 \tag{12}$$

We are now ready to evaluate the energy loss due to the emission of two majoron mediated by the Higgs boson. By implementing the result of eq. (??) in eq. (??) we obtain after integration

$$\epsilon = \frac{n_e}{\rho}\frac{120\xi}{\pi^2}\frac{T^6}{m_e^4}\zeta(6) \tag{13}$$

where $\zeta(6) = \pi^6/945 \simeq 1.017$. By comparing eq. (??) with the corresponding expression for the emissivity due to single majoron emission [?], one notices that the temperature dependence is the same. This allows us to immediately obtain a bound on the majoron-Higgs boson coupling from the result of ref. [?] on the majoron coupling to electrons : $g_{se} < 1.4 \times 10^{-13}$. By simple substitution we obtain

$$\frac{\lambda_3}{2\lambda_1} < 0.96 \tag{14}$$

Before discussing the implications of the bound obtained in eq. (??) let us recall that if the mass of ρ_L is substantially smaller than the typical temperature of the star ($\sim 10\ KeV$), then the process $\gamma e \to e\rho_L\rho_L$ will give a contribution comparable to the two majoron emission, thus further reducing the value of the bound by about a factor $\sqrt{2}$. It is also worth noting that replacement of the virtual Higgs by Z^0 gives rise, in non-singlet majoron models, to the process $\gamma e \to eJ\rho_L$. However, the derivative nature of the $Z^0 J\rho_L$ coupling leads to a further (T^2/m_e^2) suppression factor in the emissivity with respect to the two previously considered processes. This is in fact analogous to the emission of neutrino-antineutrino pairs, where the emissivity depends on T^8 as well [?,?]. In passing, let us remark that in a supernova, where much higher temperatures (and densities) are present, the coherent neutral current interaction on heavy nuclei, responsible for the neutrino trapping, may also trap majorons (and ρ_L), thus spoiling any bound on the lepton breaking VEV (or other couplings) derived from the energy radiated in the SN87 event. In this case, finite temperature effects have also to be taken into account, which may lead to the restoration of the lepton number symmetry [?].

The bound of eq. (??) affects a number of phenomenological issues. Among them is the possible majoron contribution to the rare decay $K^+ \to \pi^+ + nothing$ [?]. The Higgs boson mediated $K^+ \to \pi^+ + JJ,\ \rho_L\rho_L$ processes could indeed produce a large enhancement with respect to the standard model prediction $BR(K^+ \to \pi^+ + nothing) \sim 10^{-10}$. For instance, for $\lambda_3 \sim 0.1$ and $m_H = 20\ GeV$ the Higgs induced contribution to the decay rate would produce an enhancement of more than two orders of magnitude with respect to the three neutrino-antineutrino modes, leading to a BR at the verge of the present experimental limit. Consideration of the bound of eq. (??) limits instead this contribution to be of the same order of the standard ones.

Another interesting issue is given by the possibility that the neutral Higgs boson decays mainly in "invisible" majoron modes. This problem has been studied in some detail in ref. [?]. Implementation of the bound here obtained in the aforementioned analysis shows that the branching ratio to majorons is bounded to be less than 70% for a substantial interval of the Higgs boson mass (dashed line in fig. 1 of ref. [?]).

Finally, the astrophysical bound on the Majoron-Higgs boson coupling may affect the problem of the stability of the VEV hierarchy, generally present in non-singlet majoron models [?,?]. In fact, if we were to take face value the result of eq. (??) together with the conditions imposed, in the analysis of the VEV hierarchy, on the ratio λ_3/λ_1 by considering only the gauge contributions to the effective Higgs potential, we would find somewhat problematic to reconcile the astrophysical bound with the relation required by perturbative stability in the Gelmini-Roncadelli model, namely $\lambda_3^{(T)} = 2(4-\eta)\lambda_1^{(T)}$ [?], where $\eta \equiv 4\cos^4\theta_W/(1+2\cos^4\theta_W) \simeq 1.08$. On the other hand, eq. (??) would be barely consistent with the analogous relation for the doublet majoron model, $\lambda_3 = 2\lambda_1$ [?]. We pointed out however that the presence of a heavy top quark ($m_t > 50-60\ GeV$) may substantially modify the previous relations. The new condition for the triplet majoron is indeed given by $(1-\eta t)\lambda_3^{(T)} = 2(4-\eta)\lambda_1^{(T)}$, with $t \equiv m_t^4/m_W^4$, whereas for the doublet majoron model reduces to $(1-\eta t)\lambda_3 = 2\lambda_1$. The astrophysical bound on the coupling allows us now to constrain the range of variation of the top mass. In particular, we may conclude that the requirement of stability of the VEV hierarchy in the triplet majoron model implies $m_t \gtrsim 100\ GeV$ with a milder result for the doublet majoron model.

As a final remark, let us mention that consideration of extra-contributions to the energy loss related to the the Majoron-Higgs boson coupling, as the Bremsstrahlung production of two majorons or plasmon effects are expected, for the range of densities and temperatures

relevant for red giant stars, to give only a small correction to the bound, as can be infered from the existing analysis of energy loss through neutrino [?,?] or one majoron emission [?].

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Figure Captions

Fig. 1 Leading diagrams contributing to the process $\gamma\ e \to e\ J\ J$.