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Hidden Higgs Boson Models and Stellar Energy Loss

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ABSTRACT

We study a mechanism for the energy loss in the core of a star via the Compton–like process $\gamma e \to \phi_r^* e$ with the virtual Higgs boson ϕ_r^* turning into a pair of Goldstone bosons JJ that carry away the thermonuclear energy. Astrophysical constraint upon the coupling is derived. The bound is loose enough that the decay mode $H \to JJ$ can well be the dominating channel in the hidden Higgs models.

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For various reasons, it is sometimes desirable to extend the standard $SU_L(2)\times U_Y(1)$ electroweak gauge model with additional global symmetry which is broken spontaneously. For the example of the lepton number non–conservation, Majorana neutrino mass arises and a massless Goldstone boson, dubbed Majoron,[1,2] exists as the the remnant of the original lepton number symmetry. Another example is the Peccei-Quinn symmetry[3] which solves the strong CP problem and gives a light axion,[4] the pseudo–Goldstone boson with mass solely due to the QCD anomaly.

It has been noticed[5,6] that the physical Higgs boson H^0 can couple naturally to a pair of the Goldstone bosons J. Such an amplitude comes from the Higgs potential of the neutral component ϕ_0 of the electroweak doublet and the neutral s field which carries the quntum number of the global symmetry,

$$V(\phi,s)=-\mu^2\phi_0^\dagger\phi_0+\lambda(\phi_0^\dagger\phi_0)^2-m^2s^\dagger s+l(s^\dagger s)^2+\beta(s^\dagger s)(\phi_0^\dagger\phi_0)\ . \qquad (1)$$

This potential respects the $U(1)$ global symmetry of $s\rightarrow se^{i\delta}$. The real parts of ϕ_0 and s spontaneously develop vacuum expectation values v_2 and v_s respectively. We separate the real parts ϕ_r, s_r and imaginary parts ϕ_i, s_i of the dynamical variables as follows:

$$\phi_0=\frac{1}{\sqrt{2}}(v_2+\phi_r+i\phi_i),\quad s=\frac{1}{\sqrt{2}}(v_s+s_r+is_i)\ . \qquad (2)$$

The mass term is

$$\mathcal{L}_m=-\frac{1}{2}\begin{pmatrix}\phi_r & s_r\end{pmatrix}\begin{pmatrix}2\lambda v_2^2 & \beta v_s v_2\\ \beta v_s v_2 & 2lv_s^2\end{pmatrix}\begin{pmatrix}\phi_r\\ s_r\end{pmatrix}\ . \qquad (3)$$

For nonvanishing β, ϕ_r and s_r mix with each other. The physical mass eigenstates H and σ become

$$\begin{pmatrix}H\\ \sigma\end{pmatrix}=\begin{pmatrix}\cos\theta & \sin\theta\\ -\sin\theta & \cos\theta\end{pmatrix}\begin{pmatrix}\phi_r\\ s_r\end{pmatrix}\ . \qquad (4)$$

From the potential in Eq. (1), the induced tri–boson couplings of the Higgs boson or the σ boson to the s_i^2 are

$$\mathcal{L}_{\text{tri.}} = -\frac{1}{2}(\beta v_2 \phi_r + 2 l v_s s_r) s_i^2 \ . \tag{5}$$

The coefficients are proportional those in the bottom row of the mass matrix in Eq. (3). From the eigenvalue condition, it is straightforward to obtain

$$\mathcal{L}_{\text{tri.}} = -\frac{1}{2v_s}(\sin\theta \ m_H^2 H + \cos\theta \ m_\sigma^2 \sigma) s_i^2 \ . \tag{6}$$

If $T_3(s) = 0$, s_i is the Goldstone boson J. When $T_3(s) \neq 0$, *e.g.* s is the $T_3 = -1$ component of a triplet of lepton number $L = 2$ in the Gelmini–Roncadelli (GR) model,[2] then s_i will mix with ϕ_i. In this case, the astrophysical bound[7] requires $v_s \ll v_2 \simeq (\sqrt{2}G_F)^{-\frac{1}{2}}$ as in the GR model, the mixing can be neglected and we can still consider

$$s_i \simeq J \ . \tag{7}$$

Eq. (6) implies that the Higgs boson H and the σ boson can decay to the Goldstone boson pair JJ dominantly for a reasonable mixing angle θ.

It is interesting to put constraint on these tri–boson couplings using the low energy physics. Naively, one will take the coefficient in the Hs_i^2 term of Eq. (6) as the tri–boson coupling for all values of momentum. This will overestimate the amplitude if the σ boson contribution is not included. For an off–shell virtual ϕ_r^* of momentum p emitted by the matter field, the decay amplitude $\phi_r^* \to JJ$ suffers delicate cancellation between the H and σ contributions,

$$\begin{aligned} \mathcal{A}(\phi_r^* \to JJ) =& \frac{1}{v_s} \cos\theta \ \sin\theta \ (\frac{m_H^2}{p^2 - m_H^2} - \frac{m_\sigma^2}{p^2 - m_\sigma^2}) \ , \\ =& \frac{1}{v_s} \cos\theta \ \sin\theta \ (\frac{p^2}{p^2 - m_H^2} - \frac{p^2}{p^2 - m_\sigma^2}) \ . \end{aligned} \tag{8}$$

The amplitude vanishes at zero momentum as a property of the Goldstone boson, therefore it is difficult to obtain stringent limits upon these couplings. This can be seen more clearly if the Goldstone boson is identified as the phase of the neutral s field, instead of the imaginary part.[5] In such phase representation, the Goldstone boson appears with the derivative couplings. However, the above observation reaffirms that it is legitimate to use the formalism of Eq. (2) in perturbative analysis as long as all the fields are properly included.

In this paper, we study the constraint among the mass m_σ of the σ boson, the vacuum expectation value v_s associated with the Majoron and the mixing angle θ between the Higgs boson H and the σ boson. The constraint comes from the energy loss in the star[8] due to the Compton–like process $\gamma e \rightarrow \phi_r^* e$ of the bremsstrahlung[9] of a virtual ϕ_r boson which turns into the escaping Goldstone bosons $\phi_r^* \rightarrow (\sigma, H)^* \rightarrow JJ$ with the stellar energy carried away. We know that the single Goldstone boson bremsstrahlung[7] $\gamma e \rightarrow Je$ is suppressed by the vanishingly small Yukawa coupling $g_{e\bar{e}J}$ for the reason of either the scale hierarchy[2] or the high loop effect.[10] The mixing angle θ is an independent parameter from $g_{e\bar{e}J}$ and hence the J pair bremsstrahlung imposes new constraint on the models. Note that in the models[2,10] of our discussion with $v_s \leq v_2$, s_r does not couple to the electron.

There are also other elementary processes for the the stellar energy loss, such as the Primakoff effect and the e^+e^- annihilation. However, from the previous calculations[8,9,7] of the axion emission in the stars, the Compton–like process is the most important one for the energy loss in the red giant, which gives the strongest bound. So we will concentrate our attention to the Compton–like process.

For the star of temperature much less than $m_e c^2/k = 593 \times 10^{7}{}^\circ$K and of density much lower than $m_e/\lambda_e^3 = 1.6 \times 10^6 \mathrm{g\ cm}^{-3}$, we can use the nonrelativsitic approximations, *i.e.* the initial electron is at rest and the incident photon energy $E(\ll m_e)$ is equal to the energy loss per reaction. We also ignore the plasmon effect. The process (see Fig. 1) $\gamma e \rightarrow JJe$ is divided into two stages $\gamma e \rightarrow \phi_r^* e$ and $\phi_r^* \rightarrow (\sigma, H)^* \rightarrow JJ$ with the virtual mass m of the intermediate ϕ_r boson being integrated,

$$\sigma(\gamma e \to JJe) = \int_0^{E^2} \sigma(\gamma e \to \phi_r^* e) \frac{\theta^2}{v_s^2} \frac{dm^2}{32\pi^2} \frac{m^4}{(m^2 - m_\sigma^2)^2 + m_\sigma^2 \Gamma_\sigma^2} , \tag{9}$$

for the case $\theta \ll 1$ and $kT \ll m_H$. The cancellation of Eq. (8) has been incorporated. The subprocess cross section is[11]

$$\sigma(\gamma e \to \phi_r^* e) = \frac{\sqrt{2}}{3} \alpha G_F (1 - m^2/E^2)^{\frac{3}{2}} . \tag{10}$$

Therefore, the energy loss Q per unit time for each unit mass in the core of the star composed of nucleus of charge Z_N and mass number A_N is given by

$$\begin{aligned} Q &= \left(\frac{Z_N}{A_N m_n}\right) \frac{1}{\pi^2} \int_0^\infty \frac{E^3 \sigma(\gamma e \to JJe) dE}{\exp(E/kT) - 1} \\ &= \left(\frac{Z_N}{A_N m_n}\right) \frac{\sqrt{2} \alpha G_F \theta^2}{96\pi^4 v_s^2} T^6 \mathcal{F}\left(\frac{m_\sigma}{T}, \frac{\Gamma_\sigma}{T}\right) , \end{aligned} \tag{11}$$

with m_n as the nucleon mass. The dimensionless function $\mathcal{F}$ is defined as

$$\mathcal{F}(\beta, \gamma) = \int_0^\infty \frac{dx x^3}{e^x - 1} \int_0^{x^2} \frac{dy^2 y^4 (1 - y^2/x^2)^{\frac{3}{2}}}{(y^2 - \beta^2)^2 + \gamma^2 \beta^2} . \tag{12}$$

We can extract the γ^{-1} leading term in the narrow width approximation,

$$\mathcal{F}(\beta, \gamma) \simeq \pi \beta^3 \gamma^{-1} g(\beta) + f(\beta) . \tag{13}$$

Here the residual function $g(\beta)$ is

$$\int_\beta^\infty dx x^3 (e^x - 1)^{-1} (1 - \beta^2/x^2)^{\frac{3}{2}} , \tag{14}$$

and the remaining finite part is $f(\beta)$. Using the decay width,

$$\Gamma_\sigma = \frac{m_\sigma}{32\pi}\left(\frac{m_\sigma}{v_s}\right)^2 , \tag{15}$$

we obtain

$$Q = \left(\frac{Z_N}{A_N m_n}\right)\frac{\sqrt{2}\alpha G_F \theta^2}{3\pi^2}\left[T^4 g\left(\frac{m_\sigma}{T}\right) + \frac{1}{32\pi^2}\frac{T^6}{v_s^2} f\left(\frac{m_\sigma}{T}\right)\right] . \tag{16}$$

The g term corresponds to the on–shell σ production, and the f term is related to the off–shell (σ, H) contribution. The functions g and f are numerically given in Fig. 2. It is known that

$$g(0) = 6\zeta(4) = \pi^4/15, \quad f(0) = 48\zeta(6) = 16\pi^4/315 , \tag{17}$$

and also the asymptotic forms of g and f for large β are:

$$\begin{aligned} g(\beta) &\sim \frac{3}{4}\sqrt{\pi}(2\beta)^{\frac{3}{2}}e^{-\beta} \\ f(\beta) &\sim 18432\beta^{-4}\zeta(10) = 2048\pi^{10}\beta^{-4}/10395 . \end{aligned} \tag{18}$$

The zeta function is $\zeta(x) = \sum_{n=1}^{\infty} n^{-x}$. When $T \gg m_\sigma$, Q depends on temperature as the combination of two terms T^6 from f and T^4 from g. They can be understood as the products of the factors: the photon density ($\propto T^3$), the missing energy ($\propto T$) and the 3–body phase space ($\propto T^2$) or the 2–body phase space ($\propto 1$). On the other hand, when $m_\sigma \gg T$ at low temperature, Q depends on temperature as T^{10} because of the numerator m^4 in Eq. (9).

The energy generation rate[12] from the 3α process is $Q_{3\alpha} \simeq 100$ erg g^{-1}sec^{-1} for the helium burning red giant at a temperature 10^8 °K. The stability condition $Q < Q_{3\alpha}$ gives contraint upon θ, m_σ and v_s as illustrated in Fig. 3. For example,

$$\begin{aligned} &\theta < 3\times10^{-8} \quad \text{when} \quad m_\sigma = 0 , \\ \text{or} \quad &\theta < 5\times10^{-5} \quad \text{when} \quad m_\sigma = 200 \text{ KeV} \quad \text{and} \quad v_s = 100 \text{ KeV} . \end{aligned} \tag{19}$$

Recently, Dearborn[13] *et al.* studied the relation between the energy loss and the ignition of the stellar helium core. They were able to improve the upper bound on $g_{e\bar{e}J}$. We expect that a moderate improvement on the constraint can be obtained following their procedure.

Indirectly, the stellar energy loss gives limit upon the HJJ vertex in Eq. (6),

$$\begin{aligned}\mathcal{L}_{HJJ} &= -\frac{\kappa m_H^2}{2v_2} HJJ\ ,\\ \kappa &= \theta \frac{v_2}{v_s}\ .\end{aligned} \tag{20}$$

To require the f contribution of the virtual (σ, H) effect in Eq. (16) less than the thermoenergy generation rate $Q_{3\alpha}$, we have

$$Q_{3\alpha} > \Big(\frac{Z_N}{A_N m_n}\Big) \frac{\alpha G_F^2 \kappa^2}{48\pi^4} T^6 f\big(\frac{m_\sigma}{T}\big)\ . \tag{21}$$

The upper bound of κ is given in Fig. 4. For large m_σ, the bound is simplified with Eq. (18),

$$Q_{3\alpha} > \Big(\frac{Z_N}{A_N m_n}\Big) \frac{384\zeta(10)\alpha G_F^2 \kappa^2}{m_\sigma^4 \pi^4} T^{10}\ . \tag{22}$$

The bound becomes weaker as m_σ increases. For models[10] with v_s of the size of the electroweak scale, there is no useful bound at all.

From the above analysis, the branching fraction of $H \to JJ$ normalized to that of the lepton pair $l\bar{l}$ channel is

$$\frac{\Gamma(H \to JJ)}{\Gamma(H \to l\bar{l})} = \kappa^2 \Big(\frac{m_H}{2m_l}\Big)^2 \Big(1 - \frac{4m_l^2}{m_H^2}\Big)^{-\frac{3}{2}}\ ,$$

which is only subjected to a very weak constraint. The Higgs boson can decay dominantly into the Goldstone boson pair.

In conclusion, we show that the virtual ϕ_r Higgs boson bremsstrahlung in the core of a star is a potential source of stellar energy loss. Astrophysical constriant upon the coupling κ of the tri–boson vertex HJJ is derived. The bound is loose enough that the decay mode $H \to JJ$ can well be the dominating channel.

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FIGURE CAPTIONS

1) Digrams for the virtual ϕ_r boson bremsstrahlung $\gamma e \to JJe$.

2) The functions f and g for the off–shell (σ, H) and the on–shell σ contributions versus $\beta(= m_\sigma/T)$ in Eqs. (13) and (14).

3) The θ upper bound from the constraint due to the energy loss in red giants versus m_σ for variou values of v_s. The curve labelled with σ is the contraint due to the on–shell σ production only.

4) The upper bound of κ of the $H \to JJ$ amplitude versu m_σ.

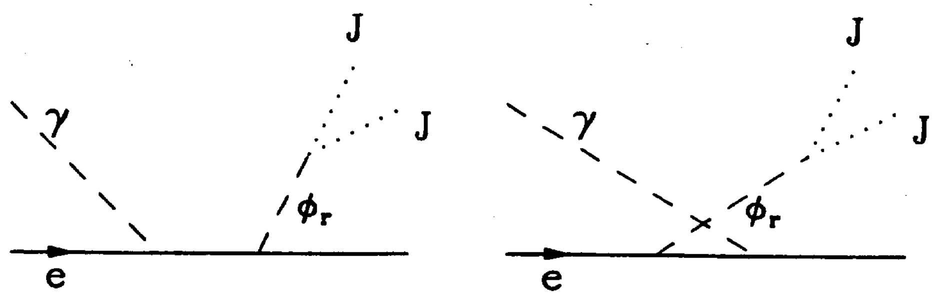

Fig. 1

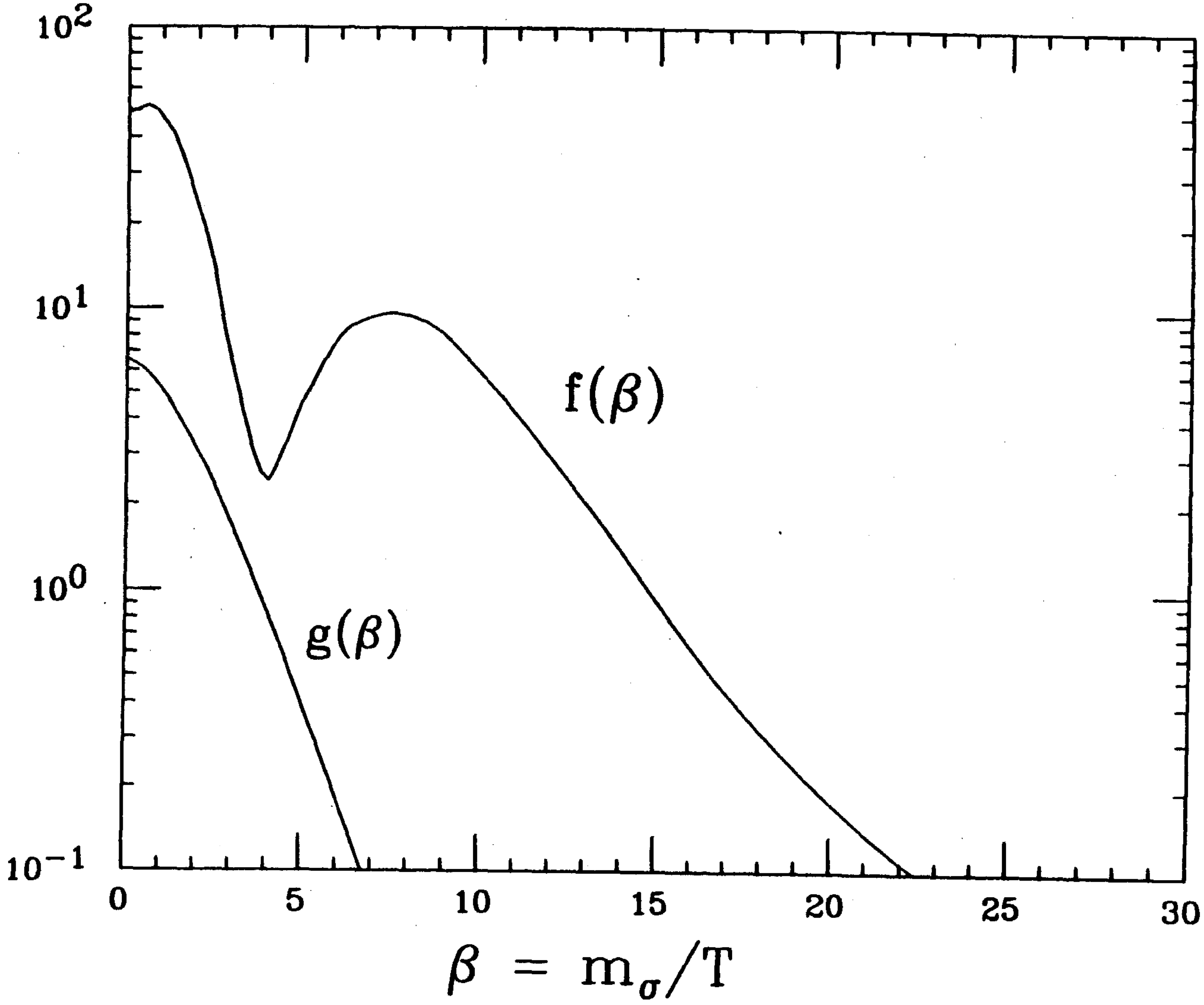

Fig. 2

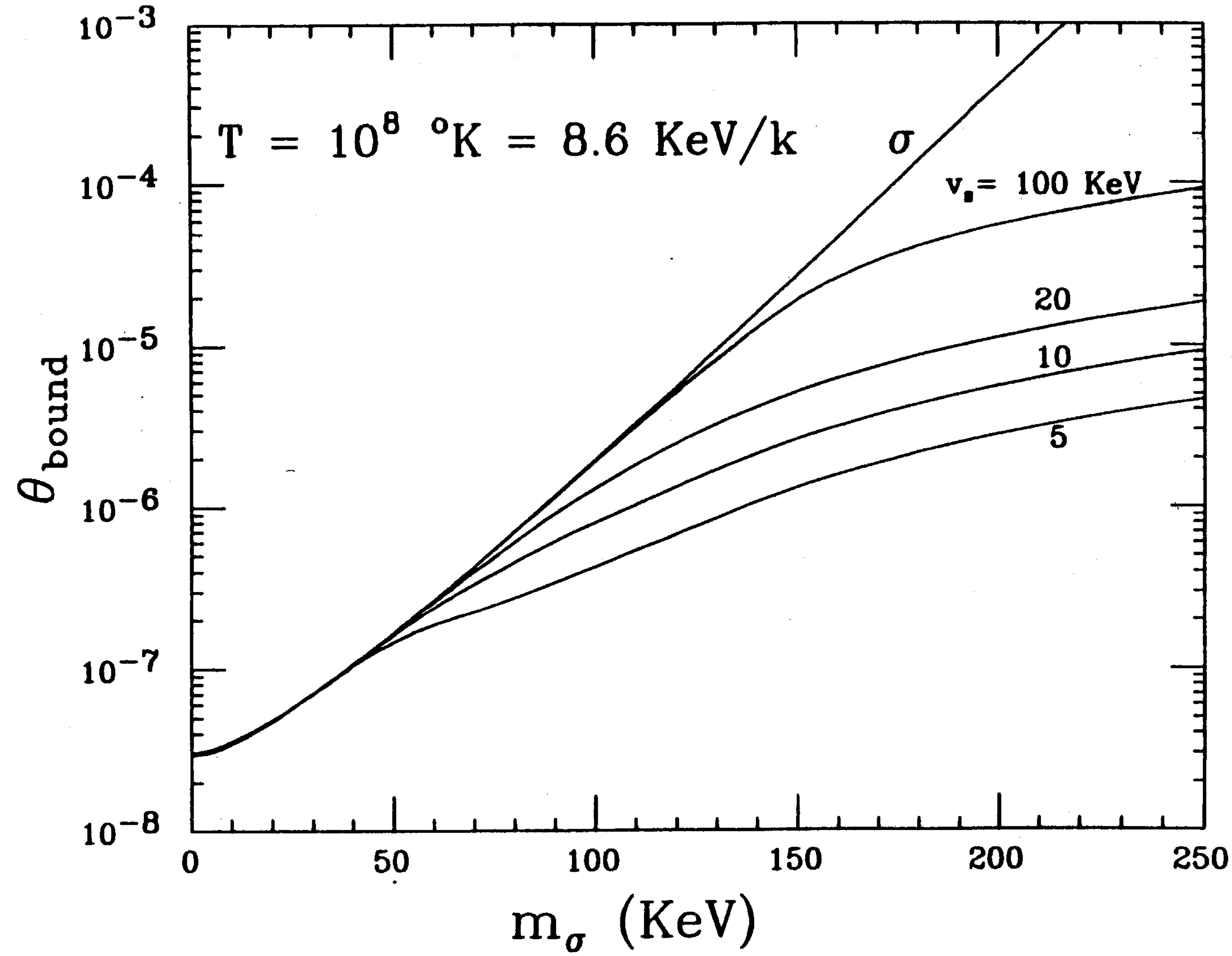
T = 10^8 °K = 8.6 KeV/k
σ
v_s = 100 KeV
20
10
5
θ_bound
m_σ (KeV)
10^-3
10^-4
10^-5
10^-6
10^-7
10^-8
0
50
100
150
200
250

Fig. 3

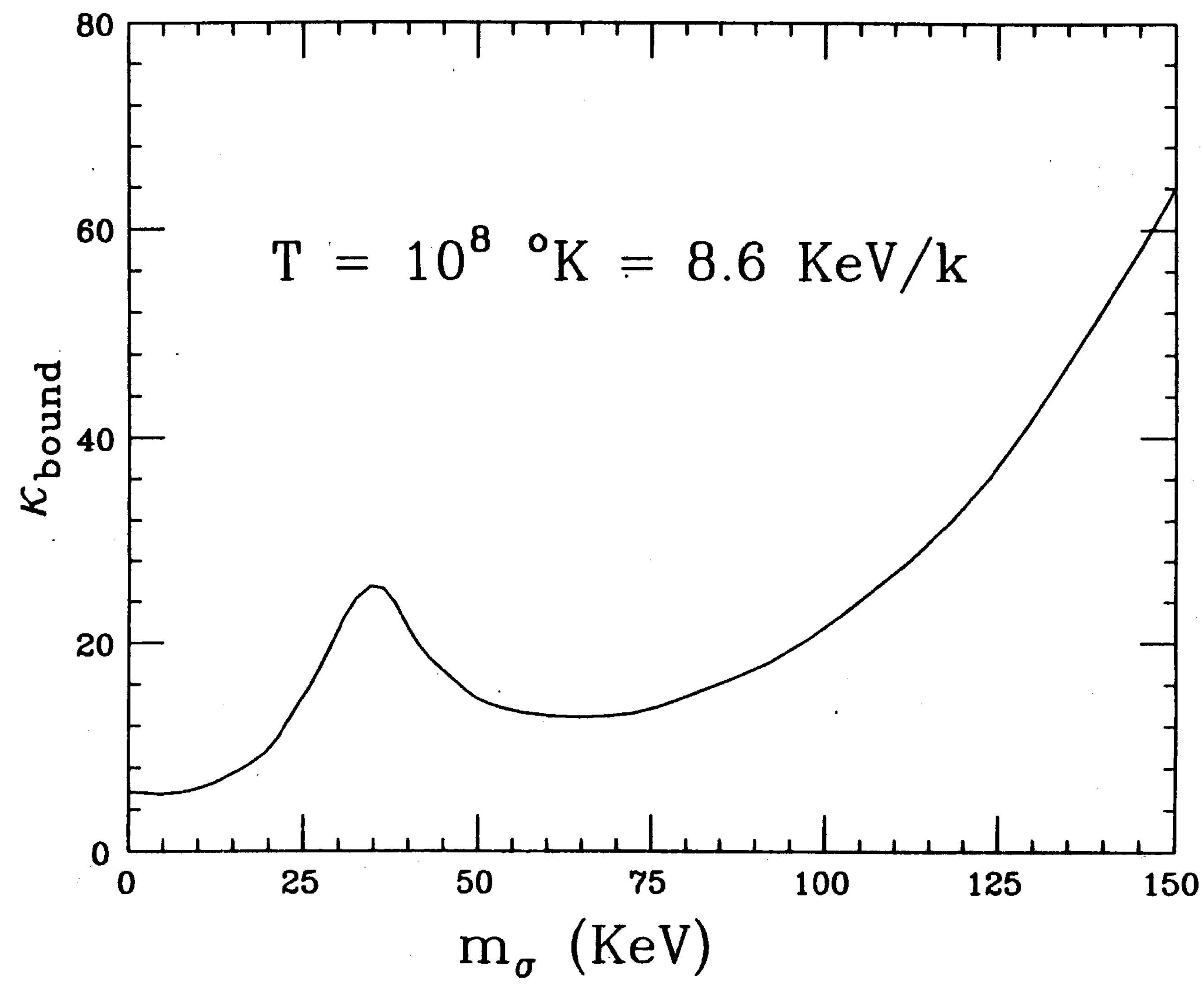
80
60
40
20
0
κ_{bound}
$T = 10^8\ ^oK = 8.6$ KeV/k
0
25
50
75
100
125
150
m_σ (KeV)

Fig. 4 .