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Astrophysical bounds on the masses of axions and Higgs particles

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Lower bounds on the mass of a light scalar (Higgs) or pseudoscalar (axion) particle are found in three ways: (1) by requiring that their effect on primordial nucleosynthesis not yield a deuterium abundance outside present experimental limits, (2) by requiring that the photons from their decay thermalize and not distort the microwave background, and (3) by requiring that their emission from helium-burning stars (red giants) not disrupt stellar evolution. The best bound is from (3); it requires the axion or Higgs-particle mass to be greater than about 0.2 MeV.

I. INTRODUCTION

It has recently been shown that astrophysical considerations may limit the properties of weakly interacting massive neutral leptons.[1-5] The arguments derive from standard big-bang cosmology: Weakly interacting neutral particles present in the early universe "freeze" out of thermal equilibrium at about 10 msec. The known upper limit on the mass density of the present universe requires the neutral particles to be unstable if their masses lie between, about, 50 eV and 5 GeV. If the decay of these leptons produces photons then stringent limits on their lifetimes can be found; this was done in Refs. 2 and 3. Here we will use the same arguments to limit the properties of light scalar or pseudoscalar particles. We can put upper bounds on the lifetimes of these particles in three ways. One is to insist that the photons from their decay not change, beyond certain limits, the calculated amount of deuterium predicted to have been produced in primordial nucleosynthesis. The second bound comes from requiring that the photons produced in the decay have time to thermalize so as not to distort the microwave blackbody background. A bound from thermalization has been previously derived by Sato and Sato[6] for (scalar) Higgs particles in a prescient paper. The third bound comes from requiring the effects on stellar evolution from axion or Higgs-particle emission not drastically reduce the lifetime of red giants. A similar bound was also calculated by Sato and Sato.[7]

There are a number of reasons to believe that the mass of the Higgs *scalar* particle is large, greater than a few GeV.[8] These arguments are, however, model dependent, and, in most cases, they assume the minimal set of Higgs particles in the usual $SU(2)\otimes U(1)$ gauge theory. For instance, in a recently proposed[9] gauge theory based on the group $SU(2)_L\otimes SU(2)_R\otimes U(1)$, not all of the scalar particles couple to fermions, some are used only to give vector gauge particles mass; this makes their detection difficult and invalidates most of the arguments restricting their mass. We believe that it would not be impossible for a very light Higgs particle to exist and to have escaped detection. It has been suggested that a light *pseudoscalar* particle (called an axion) might exist. This particle has been recently proposed[10] to explain the lack of *CP* violation in the strong interactions. A guess of the mass of the axion has been made (≈0.05–0.5 MeV) (Ref. 10) but again the estimate of its mass is uncertain, and it would not be impossible for more detailed theoretical calculations to predict it to be appreciably lighter.

In this paper we consider the possibility of the existence of a light scalar or pseudoscalr particle with interactions similar to the Higgs particle or axion. Our limits are less model dependent than limits based on current algebra or the static quark model.[11] Our best limit is from red giants and we find the mass of the axion must be greater than about 0.2 MeV.

II. LIGHT BOSONS IN THE EARLY UNIVERSE

The evolution of the universe in the presence of scalar or pseudoscalar bosons is very similar to the case of massive neutral leptons. Since these methods are well documented,[2,3] the details will be omitted. The number density of ϕ (unless otherwise noted ϕ is a scalar *or* pseudoscalar) at decoupling is calculated assuming $\phi+e\rightarrow\gamma+e$ is the process keeping the particles in equilibrium. The ϕ number density at decoupling is shown in

TABLE I. The number density of ϕ after decoupling as a function of its mass. If ϕ does not couple to electrons, n_ϕ/n_γ will be 0.5. It has been assumed that the coupling of ϕ to electrons is $\sqrt{G}\, m_e$, that is, the Higgs-particle coupling.

$\frac{m_\phi}{m_e}$	$\frac{n_\phi}{n_\gamma}$(scalar)	$\frac{n_\phi}{n_\gamma}$(pseudoscalar)
1.5	$<10^{-6}$	$<10^{-6}$
0.5	0.02	0.03
0.2	0.19	0.24
0.1	0.34	0.39
0.05	0.44	0.46
0.01	0.50	0.50

Table I as a function of the ϕ mass, assuming a Yukawa coupling to electrons with a coupling constant expected in gauge theories $\sqrt{G_F}\, m_e$.[8,10] In the usual spontaneously broken theories the axion has the same coupling to fermions as does the Higgs particle since they are combinations of the imaginary and real parts, respectively, of the original Goldstone bosons. Basically the ϕ interacts strongly enough with fermions to stay in equilibrium until the fermions annihilate. If the ϕ do not couple to electrons then they are always relativistic when they decouple so $n_\phi/n_\gamma = 0.5$. (At decoupling the ϕ look like photons but only have one spin state.)

Now consider deuterium. Its density is thought to have been almost wholly created in primordial nucleosynthesis. The predicted ^{2}H density[12] falls very sharply with increasing value of the ratio of baryons to photons at the time of nucleosynthesis. If the mass of weak scalar or pseudoscalar particles is small their lifetime will be large and then the amount of deuterium that could have been left after nucleosynthesis will be too small. This, again, is because the amount that is not converted to ^{4}He depends sensitively on the entropy per baryon.[12] We measure the entropy per baryon today (or number of photons per baryon) and, when we compare with the amount of deuterium observed, have limits on how much it could have been shifted after nucleosynthesis.[3] If the particles live too long, the entropy created by their decay into photons will shift the entropy per baryon too much. As a very good approximation[3] we may say that this happens if the ϕ live long enough to dominate the universe; that is, the bound on the ϕ lifetime may be taken as the time of domination by ϕ. By knowing the values of n_ϕ given in Table I the time of domination can be calculated as in Ref. 2; this is curve 1 of Fig. 1. Values of the lifetime τ below curve 1 are allowed. The curve goes up for large masses because the number

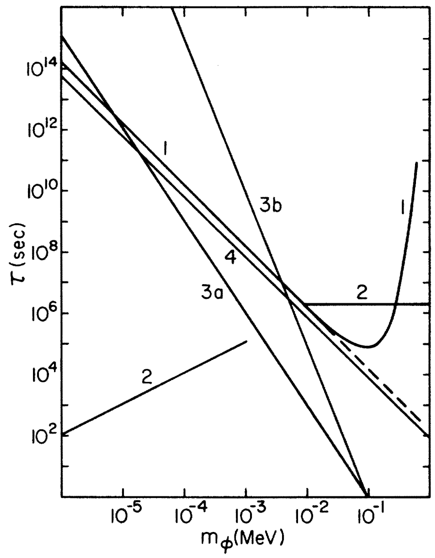

FIG. 1. Curve 1 is the maximum ϕ lifetime allowed by requiring that the decay photons not increase the entropy per baryon after nucleosynthesis. Curve 2 is the maximum ϕ lifetime allowed by requiring the decay photons thermalize. Curves 3(a) and 3(b) are the lifetimes of ϕ expected in gauge theories ($a \approx 10^{-3}$ sec MeV$^3/m_\phi{}^3$; $b \approx 0.8\times10^{-5}$ sec MeV$^5/m_\phi{}^5$). Curve 4 is the lifetime for thermalization via Thomson scattering.

denisty of ϕ drops rapidly (see Table I.) If the ϕ do not couple to electrons the number density remains large and the dashed curve is the bound.

If $n_\phi/n_\gamma < 0.03$, then no bound can be drawn from thermalization, since ϕ decay would distort the background spectrum only six (two photons from decay) parts per hundred. Therefore from Table I, if the particle couples to electrons, there is hope of a bound from thermalization only if $m < 0.5\, m_e$. If the decay photons thermalize through the reaction $\gamma + e \to \gamma + \gamma + e$, it results in the requirement that the lifetime for $\phi \to \gamma\gamma$ be less than the right-hand branch of curve 2 of Fig. 1.[2] For masses greater than 10^{-2} MeV it is possible for the decay photons to be more energetic than the background. Therefore in this case the decay photons can thermalize to become the blackbody background if the lifetime is less than about 10^6 sec. If m_ϕ is less than 10^{-3} MeV, the decay photons will not be energetic enough and it is necessary to require the lifetime to be less than the left-hand branch of curve 2 in order for the additional γ's to decrease in number and increase in energy and thus not to distort the background. In calculating the left-hand branch of curve 2 we have estimated the following effects: (1) We have taken into account stimulated absorption. This enhances the cross section for $\gamma + \gamma + e \to \gamma + e$ by about 100 (for initial γ energies one half of the final γ energy) and raises the lifetime by about 20. (2) The lifetime must be decreased by a factor of m_ϕ/kT to take into account time dilation; the photons must thermalize in the "laboratory" while the proper lifetime is measured in the rest system.

The lifetime predicted for the ϕ is less dependent on the particulars of a gauge model than its mass. (We assume $m_\phi < 2m_e$ so $\phi \to e^+e^-$ is not allowed.) Most models predict something around $10^{-3}/m_\phi^3$ sec where m_ϕ is in MeV. This is shown as curve 3(a) of Figure 1. For this mass-lifetime relation, the lower bound on the mass from nucleosynthesis is not very restrictive (~5 eV). On the other hand, a particle of mass less than ~2 $\times 10^{-3}$ MeV would live too long for its decay products to thermalize. This result is somewhat more restrictive than the previous upper bound of 10^{-4} MeV based on similar thermalization arguments.[6] For other mass-lifetime relations curve 3(a) can easily be scaled up or down; however, it is so steep that the bounds from thermalization will not change significantly. For example, curve 3(b) is an estimate of the axion lifetime that takes into consideration the possibility of ϕ-π^0 mixing,[13] $\tau = 8\times 10^{-5} m_\phi^{-5}$ (MeV). The thermalization bound remains about the same, but now the nucleosynthesis bound is comparable.

One exception tothe above thermalization bounds should be noted. Thomson scattering proceeds at a very much faster rate than either $\gamma + e \to \gamma + \gamma + e$ or $\gamma + \gamma + e \to \gamma + e$, so that a present thermal background could result from Thomson scattering alone. The energy density of the $\gamma - \phi$ system before decay is

$$\rho = aT^4 + \frac{n}{2} m_\phi + \frac{3}{2}\frac{n}{2} kT \,. \tag{1}$$

The number of photons after decay is larger than before decay by a factor of 2 ($\phi \to 2\gamma; n_\phi = n_\gamma/2$),

$$n' = 2n \,. \tag{2}$$

The energy density after decay is

$$\rho = aT'^4 \,. \tag{3}$$

Since in a blackbody distributionthe number of density is related to the temperature as

$$n = \frac{0.37aT^3}{k} \,, \tag{4}$$

(2) implies the final temperature will be

$$T' = 2^{1/3} T \,. \tag{5}$$

Using (2), (4), and (5) in equating (1) and (3) gives

$$kT = 0.15 m_\phi \,. \tag{6}$$

If the time at decay is $2\times 10^{20}/T^2$ sec, and if the mass is related to the lifetime as $t = 10^{-3}/m_\phi^3$, there is a unique mass that may be thermalized by scattering: $m_\phi = 1.5\times 10^{-5}$ MeV. For any mass-lifetime relation Thomson scattering will thermalize the photons if the lifetime is given by curve 4. The small-mass section of curves 2 and 4 are not affected by time dilation since in both cases the decay time is proportional to m_ϕ^{-2}.

In summary, information from cosmology indicates that an unstable scalar or pseudoscalar Higgs particle should either have a mass greater than 10^{-3} MeV, or else have a lifetime given by curve 4 of Fig. 1.

III. LIGHT BOSONS AND STELLAR EVOLUTION

In this section we consider the effect of light bosons on stellar evolution. After the depletion of hydrogen, stars of a few solar masses enter a stable configuration where helium is burned in the core.[14] A core of about 0.5 $M_\odot$ is present prior to helium ignition in a 5 $M_\odot$ star. The helium starts to burn at a temperature of about 10^8 °K and and a density of about 10^4 g cm^{-3}.[14,15] Stars in this stage form red giants. Red giants are observed and their properties are in agreement with evolutionary calculations.

Sato and Sato[7] considered the effect of Higgs scalars in stellar evolution and found that the mass of the Higgs particle should be greater than 0.35 MeV if present models of stellar evolution are not to be affected. When they performed the calculation the Higgs particle was believed to be greater than a few GeV; now with the proposal of a light pseudoscalar in the mass range 0.1 to 0.5 MeV, their approximate observations warrant expansion and improvement.

We will calculate the energy loss due to axion emission in a 0.5 $M_\odot$ helium core for two processes. The largest of the two processes gives limits on the mass of the axion, one assuming present evolutionary models of helium burning stars, the other limit independent of any such model.

The first process considered is the Primakoff process,[16] $\gamma + Z \to \phi + Z$, shown in Fig. 2. The cross section for this process near threshold is

$$|v|\sigma = 64\pi\alpha Z^2 \frac{\omega\Gamma(\phi\to 2\gamma)}{m_\phi^2} \frac{(\omega^2 - m_\phi^2)^{1/2}(\omega - m_\phi)}{(m_\phi^2 - 2\omega m_\phi)^2} \,, \tag{7}$$

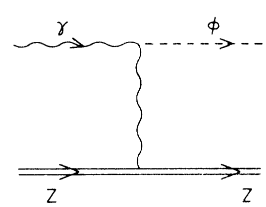

FIG. 2. $\gamma + Z \to \phi + Z$ via the Primakoff process.

TABLE II. Energy loss due to axion emission from a 0.5 $M_\odot$ helium core at 10^4 g cm^{-3} and 10^8 °K. The 3α process at this temperature and density produces 10^2 erg g^{-1}sec^{-1}. A and B are the two possibilities for $\Gamma(\phi\to 2\gamma)$ discussed for the Primakoff process. Also given is the loss rate for $\gamma+e\to\phi+e$ both corrected and uncorrected for reabsorption and decay.

	Q (erg g^{-1}sec^{-1})			
	Primakoff (10)		Compton (11)	
m_ϕ	$\gamma+Z\to\phi+Z$		$\gamma+e\to\phi+e$	
(MeV)	A	B	uncorrected	corrected
10^{-3}	...	...	1.1×10^{15}	1.1×10^{15}
10^{-2}	...	...	1.0×10^{15}	3.4×10^{14}
0.1	6.1×10^{8}	7.6×10^{8}	5.9×10^{12}	1.4×10^{11}
0.2	1.5×10^{4}	7.3×10^{4}	6.6×10^{8}	6.8×10^{5}
0.25	6.1×10^{1}	4.8×10^{2}	4.4×10^{6}	3.6×10^{3}
0.3	2.4×10^{-1}	2.7×10^{2}	2.5×10^{4}	9.4×10^{-1}
0.35	<1	<1	1.3×10^{2}	<1

where ω is the energy of the initial photon. The energy loss (uncorrected for reabsorption or ϕ decay) is found by integrating (7), times the energy of ϕ ($\approx\omega$), over the phase space of the initial particles

$$Q=\int dn_Z\int dn_\gamma E_\phi\sigma|v|\,. \tag{8}$$

The energy loss in erg g^{-1} sec^{-1} becomes

$$Q=\frac{64\alpha Z^2}{\pi\rho}\frac{\Gamma(\phi\to2\gamma)}{{m_\phi}^2}n_Z(kT)^3\times\int_y^\infty dx\,\frac{x^4(x^2-y^2)^{1/2}(xy)}{(e^x-1)(y^2-2xy)^2} \tag{9}$$

where $y=m_\phi/kT$, n_Z is the number density of ^{4}He, and ρ is the mass density. In Table II we consider two possibilities for Γ: Case A has been used in the previous section,

$$\Gamma=10^3\left(\frac{m_\phi}{\text{MeV}}\right)^3\text{ sec}^{-1}\,. \tag{10a}$$

Case B uses current algebra to find the $\phi-\pi$ mixing angle[13]

$$\Gamma=1.25\times10^5\left(\frac{m_\phi}{\text{MeV}}\right)^5\text{ sec}^{-1}\,. \tag{10b}$$

The Primakoff process does not give a better bound than $\gamma+e\to\phi+e$, but it is nevertheless important since it is conceivable that the axion does not couple to electrons. In order that the energy loss from axion emission be less than the nuclear energy generation rate ($\leq10^2$ erg g^{-1} sec^{-1}) at 10^8 °K and 10^4 g cm^{-3} the mass of the axion must be greater than about 0.2 MeV.

Now consider Sato and Sato's processes, $\gamma+e\to\phi+e$. The energy loss for this processes (again uncorrected for reabsorption) is given again by (8); however, now the cross section is a complicated function of the energies and angles. Since we will be considering photon energies comparable to the electron mass, it is not possible to assume the electrons are at rest, and their phase space must be integrated over also,

$$Q=\frac{4}{(2\pi)^6}\int d^3p_1\int d^3k[\exp(E_1/kT)+1]^{-1}\times[\exp(\omega/kT)-1]^{-1}E_\phi|v|\sigma(p_1,k)\,. \tag{11}$$

The energy loss must be corrected for reabsorption. Once the axion is created, if the mean free path for reabsorption, $\phi+e\to\gamma+e$, is less than the radius of the star, not all of the axions will escape. We have approximated the suppression in the energy loss due to axion reabsorption ($\phi+e\to\gamma+e$) by calculating the cross section σ_R for reabsorption, finding the mean free path

$$\langle l\rangle=\langle n_e\sigma_R\rangle^{-1}\,, \tag{12}$$

and multiplying (11) by $(l/R_c)^2$ where R_c is the radius of the core. This is the random-walk approximation of diffusion. We find for a 0.2-MeV axion that $\langle l\rangle\sim R_c/10$ for $\rho=10^4$ g cm^{-3} and $\langle l\rangle\sim R_c/30$ for $\rho=10^5$ g cm^{-3}. If the mass of the axion is larger (smaller), $\langle l\rangle$, is somewhat smaller (larger). Treating reabsorption as a random walk is an approximation; however, we have considered other treatments of reabsorption and they do not change the limits on the axion mass. The axion energy production (11) is always *many* orders of magnitude greater than the nuclear energy generation (e.g., Table II) if the Boltzmann factor is not small; the Boltzmann factor completely determines when the axion emission is small enough.

We have also considered the possibility that the axion will decay before leaving the star. The decay length as a function of the axion energy was calculated,

$$\langle d\rangle = \gamma\beta\tau c\ , \tag{13}$$

and (11) was multiplied by $e^{-R_c/\langle d\rangle}$. For a 0.1- or 0.2-MeV axion $\langle d\rangle \sim (1-2)R_c$, and for $m_\phi < 0.1$ MeV $\langle d\rangle > R_c$. A lifetime of $10^{-3}(1\ \text{MeV}/m_\phi)^3$ sec was used in (13). Again the exact form of τ is not important and for the same reason. It is also easy to show that the axion will have enough kinetic energy to escape from the gravitational attraction of the core even if the temperature is as low as 10^6 °K.

Equation (11) was calculated numerically and is given in Table II (corrected for reabsorption and ϕ decay) for $\rho = 10^4$ g cm^{-3} and $T = 10^8$ °K. Sato and Sato have approximated (11) and their results are in good agreement with our exact (up to a three-dimensional numerical integral) results in regions where their approximations are valid. The energy generation due to helium burning at this temperature and density is 10^2 erg g^{-1} sec^{-1}.[14,15] An axion mass less than about 0.3 MeV will therefore upset models of stellar evolution.

A similar calculation may be done for more massive stars. We have calculated the energy loss for a 15 $M_\odot$ star[17] (about a 4 $M_\odot$ helium core). These stars burn helium at a slightly higher temperature, $T = 1.7 \times 10^8$ °K, and at a lower density, $\rho = 1.14 \times 10^3$ g cm^{-3}. The higher temperature results in a higher axion emission rate since the Boltzmann suppression is not as great. Applying the restriction that the energy loss be less than the energy generation at the above temperature and density configuration requires the axion's mass to be greater than 0.4 MeV. Sato and Sato consider similar conditions and require the mass be greater than 0.35 MeV.

The above arguments depend on models of stellar evolution to determine the temperature and density regions where the star should burn helium. However, if light bosons exist they may in fact *determine* the evolution of the star after hydrogen depletion and old models cannot be trusted to give the values of ρ and T.

But it is possible to consider restrictions on the axion mass that are independent of models of stellar evolution. In Fig. 3 we show configurations in the ρ, T plane where helium may burn, i.e., above the lines for a given axion mass the 3α helium burning energy generation is greater than the axion emission and helium can burn, below the lines the axion energy losses dominate the energy generation and the star is unstable. Also shown in Fig. 3 is the energy generation for helium burning as a function of temperature and density.[18]

If we require helium to burn above the 0.2-MeV line in a nondegenerate configuration (degenerate

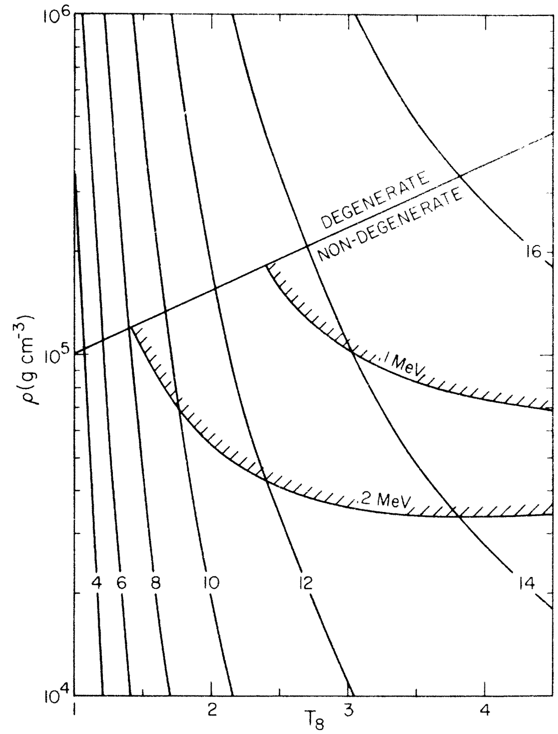

FIG. 3. The shaded areas for a given axion mass for the areas where helium will burn, that is, where the energy generation from helium burning exceeds the energy loss due to axion emission. The numbered lines are the logarithm of the 3α energy generation as a function of density and temperature; e.g., 10 is the locus of points where the 3α process produces 10^{10} erg g^{-1}sec^{-1}.

helium ignition is explosive) the presence of the axion causes the helium to burn at an extremely high rate ($\geq 10^8$ erg g^{-1} sec^{-1}) and the star will have an extremely short lifetime ($\leq$one year) in conflict with observation. If the axion is 0.1 MeV the situation is worse, but if the mass of the axion is 0.3 MeV or higher, helium can burn at a reasonably slow rate. If $m_\phi = 0.3$ MeV helium may burn at 10^8 °K and 10^4 g cm^{-3}, see Table II.

If the axion mass is as large as a few MeV it will still have a very important effect on later stages of stellar evolution such as carbon burning or supernovas.[19]

In conclusion, we have shown that present models of stellar evolution suggest the mass of the axion is greater than 0.4 MeV, and we can require the axion to be more massive than 0.2 MeV independent of evolutionary models.

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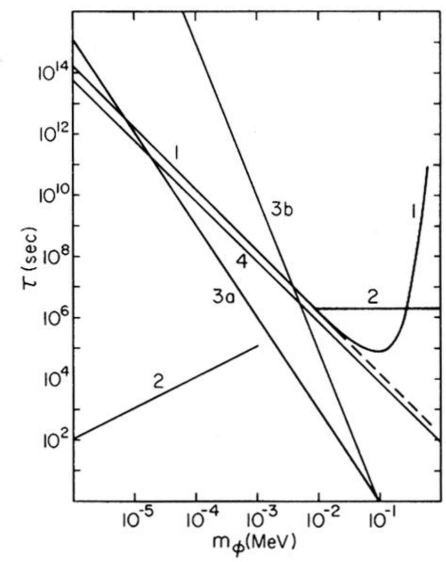

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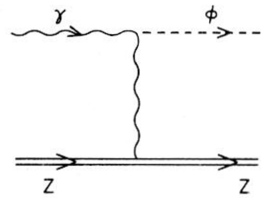
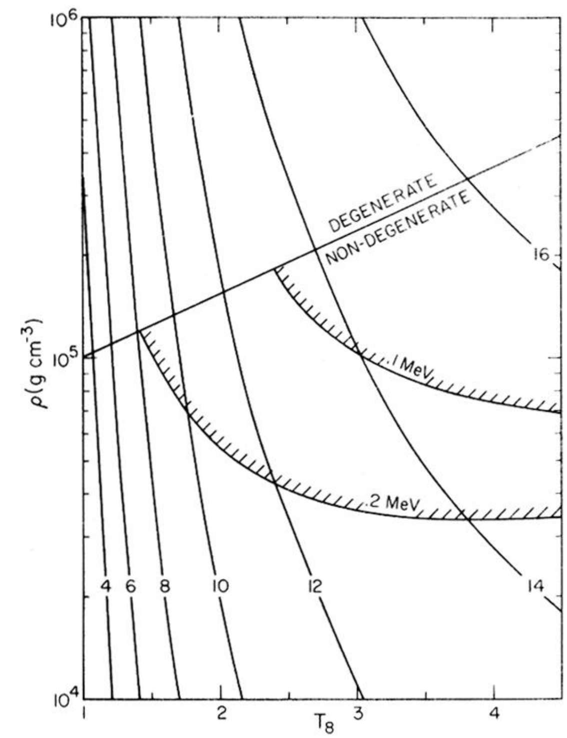


FIG. 2. $\gamma + Z \to \phi + Z$ via the Primakoff process.



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