PRAMANA
—journal of
physics

Vol. 51, Nos 1 & 2,
July/August 1998
pp. 135–143

Astrophysical constraints on superlight gravitinos

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Abstract. I review the constraints on the mass of gravitinos that follow from considerations on energy loss in stars and from big bang nucleosynthesis arguments.

Keywords. Gravitinos; stellar energy loss; big bang nucleosynthesis.

PACS Nos 14.80; 04.65; 97.60

1. Stellar energy loss [1]

1.1 *The sun*

Thermal stars like our sun are systems where thermal pressure balances gravity. An important feature of these stars is their negative specific heat. They heat up when the system looses energy. Indeed, when the star gives off energy, its internal plus gravitational energy goes down and as a consequence the star contracts. This is so because the gravitational energy itself diminishes, as required by the virial theorem. Contraction, in turn, implies a rise in temperature and extra nuclear burning because the internal energy has to compensate for the decrease of gravitational energy (virial theorem, again). Now, the added nuclear energy supply leads to an opposite cycle of expansion and cooling (negative specific heat) and hence to a lesser nuclear burning. This is a self-regulating mechanism where equilibrium is maintained by balancing the energy loss with nuclear fusion.

Weakly interacting particles, if not trapped in the interior of stars, drain energy. Hence they lead to accelerated consumption of nuclear fuel. Models of stellar evolution tolerate non-standard exotic energy drain mechanisms as long as their associated luminosity does not exceed the solar luminosity, i.e. $L \leq 10^{33}$ erg/s.

1.2 *Degenerate stars*

These are systems dominated by the Fermi pressure of electrons (white dwarfs) or of nucleons (neutron stars). In this case the internal energy (Fermi energy) is nearly independent of the temperature and a loss of energy is at the expense of gravitational contraction (which increases the internal Fermi energy but not the temperature). In fact the system actually cools.

Here also, new hypothetical particles provide for extra energy depletion which should not exceed 5×10^{-3} erg/g/s.

1.3 *Type II supernovae*

Very massive stars (much heavier than the sun) ignite all nuclear fuels up to iron. Beyond this point no more nuclear energy is available. When the iron core reaches its Chandrasekhar limit, the star enters an unstable catastrophical regime: gravitational collapse follows.

1.3.1 *Gravitational collapse*: After the thermonuclear cycle is completed, Si having transformed into Fe, collapse is triggered by the photodissociation of Fe^{56} and electron capture on nuclei. The core loses Fermi pressure from the electrons which no longer hold gravitational pressure and the core implodes. Hence, ν_e emission dominates during infall. As the core collapses temperature and density rise. For densities above $2 \times 10^{11}\,\mathrm{g/cm^3}$ neutrinos become trapped in the core. When the inner core reaches nuclear densities ($\sim 3 \times 10^{14}\,\mathrm{g/cm^3}$) a nuclear phase transition occurs, i.e. bound nuclei become free nucleons. This has a dramatic effect. The core stiffens because now non-relativistic nucleons dominate the pressure (before, the pressure of relativistic electrons and electron neutrinos could not halt collapse). The inner core bounces generating a shock wave which is responsible for the ejection of the mantle of the star.

The hydrodynamical core collapse happens in less than a second. The relevant dynamic time scales are on the order of 10 to 100 miliseconds (almost free-fall). For instance, the initial electron–neutrino burst lasts for about 5 ms. Finally, the proto-neutron star originated that way cools off to a cold neutron star in a time scale of several seconds.

Independently of the details of collapse, to form a neutron star

$$E_B \sim GM^2/R \sim 3 \times 10^{53}\,\mathrm{erg} \tag{1}$$

have beeen released. The total luminosity in electromagnetic radiation plus kinetic energy of the ejecta in a supernova (SN) explosion is $\leq 10^{51}$ erg. Emission in gravitational radiation is at most 1%. The bulk of the binding energy ($\geq 99\%$) is emitted in form of neutrinos. In the neutronization burst $\sim 10^{52}$ erg (10% of the total energy) are emitted. The rest of the energy is radiated essentially as thermal neutrinos.

Observational data from supernova 1987A (IMB, Kamiokande) [2] imply

$$E_\nu \geq 2 \times 10^{53}\,\mathrm{erg} \tag{2}$$

emitted over a diffusion period of the order of ten seconds. Comparing equations (2) and (1) we see that any additional energy drain is allowed whenever

$$L_X \leq 10^{52}\,\mathrm{erg/s}. \tag{3}$$

2. A case study: Bounds on gravitino masses

2.1 *The gravitino*

The gravitino is the spin 3/2 partner of the graviton. It is a Majorana particle with only transverse degrees of freedom before supersymmetry breakdown. After eating the goldstino, it acquires the longitudinal degrees of freedom, i.e. the $\pm 1/2$ helicities. In some recent models the gravitino can be superlight. Indeed, models where gauge interactions mediate the breakdown of supersymmetry [3], models where an anomalous U(1) gauge symmetry induces SUSY breaking [4], and no-scale models are all examples of models where a superlight gravitino can be accommodated [5]. In all of them, the gravitino is the LSP and, furthermore, its couplings to matter and radiation are inversely proportional to its mass. For very small gravitino masses, the longitudinal components of the gravitino dominate the interactions with matter so that effectively,

$$G_\mu \cong \sqrt{\frac{2}{3}} m_{3/2}^{-1} \partial_\mu \chi. \tag{4}$$

In a subclass of the models mentioned above, the scalar and pseudoscalar partners of the goldstino are also ultralight. Their couplings do also show the enhancing $m_{3/2}^{-1}$ factor [6].

2.2 *Astrophysics* [7]

2.2.1 *Supernova constraints* [8–11]: According to Luty and Ponton [12], photons have an effective coupling to gravitinos with the following structure

$$\delta L = \frac{e}{2}(M/\Lambda^2)^2 \partial^\mu \chi \sigma^\nu \bar{\chi} F_{\mu\nu} + \text{h.c.} \tag{5}$$

where M is a model dependent supersymmetric mass parameter and Λ is the supersymmetry breaking scale. This interaction provides the leading contribution to gravitino pair emissivity via electron–positron annihilation, nucleon–nucleon bremsstrahlung, plasmon decay, and photon–electron scattering followed by radiation of the gravitino pair.

The energy-loss rate (per unit volume) via $pn \to pn\tilde{G}\tilde{G}$ is,

$$Q = \int \frac{\mathrm{d}^3 k_1}{(2\pi)^3 2k_1^0} \frac{\mathrm{d}^3 k_2}{(2\pi)^3 2k_2^0} \prod_{i=1}^{4} \frac{\mathrm{d}^3 p_i}{(2\pi)^3 2p_i^0} f_1 f_2 (1-f_3)(1-f_4) \times (2\pi)^4 \delta^4(P_f - P_i) \sum_{\text{spins}} |M_{fi}|^2 (k_1^0 + k_2^0), \tag{6}$$

where $(p^0, \vec{p})_i$ are the 4-momenta of the initial and final state nucleons, $(k^0, \vec{k})_{1,2}$ are the 4-momenta of the gravitinos and $f_{1,2}$ are the Fermi–Dirac distribution functions for the initial proton and neutron and $(1 - f_{3,4})$ are the final Pauli blocking factors for the final proton and neutron. The squared matrix element can be factorized as follows,

$$\sum_{\text{spins}} |M_{fi}|^2 = (2\pi)^2 \alpha^2 (M/\Lambda^2)^4 N_{\mu\nu} G_{3/2}^{\mu\nu}, \tag{7}$$

where $N_{\mu\nu}$ is the nuclear (OPE) tensor and $G_{3/2}^{\mu\nu}$ is the gravitino tensor in the matrix element squared. The factor $N_{\mu\nu}$ is common to any bremsstrahlung process involving nucleons. It appears, e.g., in neutrino bremsstrahlung calculations and in axion bremsstrahlung calculations, and is given explicitly in reference [1]. On the other hand, $G_{3/2}^{\mu\nu}$ is a tensor specific to gravitino bremsstrahlung. It reads,

$$G_{3/2}^{\mu\nu} = k_1^\mu k_2^\nu + k_2^\mu k_1^\nu - k_1 \cdot k_2 g^{\mu\nu}. \tag{8}$$

The integration of $N_{\mu\nu}$ over the phase-space of the nucleons can be performed explicitly and the details can be found in Raffelt's book [1]. When we contract the result with the gravitino tensor $G_{3/2}$ and perform the integrals over gravitino momenta to complete the energy depletion rate, we are led to the following emissivity:

$$Q^{\mathrm{ND}}_{\mathrm{brems}} = (8192/385\pi^{3/2})\alpha^2\alpha_\pi^2(M/\Lambda^2)^4 Y_e n_B^2 T^{11/2}/m_p^{5/2} \quad (9)$$

for non-degenerate and non-relativistic nucleons (α_π is the pionic fine-structure constant, n_B is the number density of baryons, and Y_e is the mass fraction of protons). However, nucleons are moderately degenerate in the SN core. The emissivity in the (extreme) degenerate case is calculated to be,

$$Q^{\mathrm{D}}_{\mathrm{brems}} = (164\pi^3/4725)\alpha^2\alpha_\pi^2(M/\Lambda^2)^4 p_{\mathrm{F}} T^8 \quad (10)$$

with p_{F}, the Fermi momentum of the nucleons. Numerically, for the actual conditions of the star, both emissivities differ by less than an order of magnitude (about a factor of three). Since the actual emissivity interpolates between these two values, we shall adopt the smallest of the two (i.e. $Q^{\mathrm{ND}}_{\mathrm{brems}}$) to make our (conservative) estimates. We turn next to the annihilation process.

The energy loss for the process $e^+(p_1) + e^-(p_2) \to \tilde{G}(k_1) + \tilde{G}(k_2)$ can be calculated along similar lines as above. The spin averaged matrix element squared is in this case,

$$\sum_{\mathrm{spins}} |M_{fi}|^2 = (2\pi)^2\alpha^2(M/\Lambda^2)^4 E_{\mu\nu}(p_1,p_2) G^{\mu\nu}_{3/2}(k_1,k_2), \quad (11)$$

where $E_{\mu\nu}(p_1,p_2)$ equals formally the tensor $G^{\mu\nu}_{3/2}$ in eq.(8) with k_1,k_2 replaced by p_1,p_2. The luminosity then is found to be,

$$Q_{\mathrm{ann}} = 8\alpha^2(M/\Lambda^2)^4 T^4 \mathrm{e}^{-\mu/T}\mu^5 b(\mu/T)/15\pi^3 \quad (12)$$

with $b(y) \equiv (5/6)e^y y^{-5}(F_5^+F_4^- + F_4^+F_5^-)$ where $F_m^\pm(y) = \int_0^\infty \mathrm{d}x x^{m-1}/(1+\mathrm{e}^{x\pm y})$ (μ is the chemical potential of the electrons). The function $b(y) \to 1$ in the degenerate limit. Finally, our estimate of the plasmon decay luminosity is,

$$Q_{\mathrm{P}} = 16\zeta(3)\alpha^4 T^3\mu^6(M/\Lambda^2)^4/81\pi^5 \quad (13)$$

(where only transverse plasmons have been taken into account).

The process $\gamma e \to e\tilde{G}\tilde{G}$ has not been evaluated analytically but numerically has been seen to be of the same order as $e^+e^- \to \tilde{G}\tilde{G}$ [11]. Taken at face value, the bremsstrahlung rate is the largest of the four. However, Q_{brems} is overestimated since we did not consider multiple scattering effects which are present in a dense medium [1]. Indeed, as for the axion case [1], the gravitino bremsstrahlung rate probably saturates around 20% nuclear density and this should be taken into account when evaluating eq. (11). If we use now the values $T = 50\,\mathrm{MeV}$, $\mu = 300\,\mathrm{MeV}$, and $Y_e = 0.3$, eqs (11) (with $n_B \sim 0.2 n_{\mathrm{nuc}}$), (14) and (15) give

$$Q_{\mathrm{ann}} : Q_{\mathrm{brems}} : Q_{\mathrm{P}} \approx 1.2\times10^3 : 3\times10^2 : 1 \quad (14)$$

Therefore, a limit on Λ will follow from the requirement that $L_{3/2} \approx VQ_{\mathrm{ann}}$ (V is the volume of the stellar core) should not exceed 10^{52} ergs/s. This constraint on the gravitino luminosity $L_{3/2}$ implies, in turn,

$$\Lambda \geq 300\,\mathrm{GeV}(M/43\,\mathrm{GeV})^{1/2}(T/50\,\mathrm{MeV})^{11/16}(R_c/10\,\mathrm{Km})^{3/8} \quad (15)$$

or, using $m_{3/2} = 2.5 \times 10^{-4}\,\mathrm{eV}(\Lambda/1\,\mathrm{TeV})^2$,

$$m_{3/2} \geq 2.3 \times 10^{-5}\,\mathrm{eV}. \tag{16}$$

Of course, the previous calculation makes sense only if gravitinos, once produced, stream freely out of the star without rescattering. That they actually do so, for $\Lambda \geq 300\,\mathrm{GeV}$, can be easily checked by considering their mean-free-path (mfp) in the core. The main source of opacity for gravitinos is the elastic scattering off the Coulomb field of the protons:

$$\lambda = 1/\sigma n = (4/\pi\alpha^2) Y_e^{-1} \rho^{-1} m_p^{-1} (\Lambda^2/M)^4. \tag{17}$$

The thermally averaged cross-section for elastic gravitino scattering on electrons is roughly a factor $T\mu/m_p^2$ smaller than that on protons and thus it does not contribute appreciably to the opacity. Putting numbers in eq. (19) we find:

$$\lambda \simeq 1.4 \times 10^7\,\mathrm{cm}(43\,\mathrm{GeV}/M)^4(\Lambda/300\,\mathrm{GeV})^8. \tag{18}$$

On the other hand, the calculation of Q breaks down for $\lambda \leq 10\,\mathrm{Km}$, i.e. for $\Lambda \leq 220\,\mathrm{GeV}$, when gravitinos are trapped in the SN core. In this case, gravitinos diffuse out of the dense stellar interior and are thermally radiated from a gravitino-sphere $R_{3/2}$. Because in this instance the luminosity is proportional to T^4, only for a sufficiently large $R_{3/2}$ (where the temperature is correspondingly lower), the emitted power will fall again below the nominal $10^{52}\,\mathrm{erg/s}$. Consequently, gravitino emission will be energetically possible, if Λ is small enough. The gravitino-sphere radius can be computed from the requirement that the optical depth

$$\tau = \int_R^\infty \mathrm{d}r/\lambda(r) \tag{19}$$

be equal to 2/3 at $R = R_{3/2}$. Here, $\lambda(r)$ is given in eq.(19) with the density profile ansatz:

$$\rho(r) = \rho_c (R_c/r)^m \tag{20}$$

with $\rho_c = 8 \times 10^{14}\,\mathrm{g/cm^3}$, $R_c = 10\,\mathrm{Km}$ and $m = 5 - 7$ and which satisfactorily parameterizes the basic properties of SN1987A [13]. An explicit calculation renders:

$$R_{3/2} = R_c[(8Y_e/3\pi\alpha^2)(\Lambda^2/M)^4(m-1)/\rho_c R_c m_p]^{1/1-m}. \tag{21}$$

Stefan–Boltzmann's law implies for the ratio of gravitino to neutrino luminosities,

$$L_{3/2}/L_\nu = (R_{3/2}/R_\nu)^2[T(R_{3/2})/T(R_\nu)]^4, \tag{22}$$

where R_ν is the radius of the neutrinosphere. To proceed further we use the temperature profile:

$$T = T_c(R_c/r)^{m/3} \tag{23}$$

which is a consequence of eq. (20) and the assumption of local thermal equilibrium. Now, taking $m = 7$ [14], we obtain

$$L_{3/2}/L_\nu = (R_\nu/R_c)^{22/3}[(16Y_e/\pi\alpha^2)(\Lambda^2/M)^4/\rho_c R_c m_p]^{11/9}. \tag{24}$$

By demanding that $L_{3/2} \leq 0.1 L_\nu$ and using $R_\nu \simeq 30\,\mathrm{Km}$, we get

$$\Lambda \leq 70\,\mathrm{GeV}. \tag{25}$$

This in turn implies $m_{3/2} \leq 10^{-6}\,\mathrm{eV}$. Since, on the other hand, the anomalous magnetic moment of the muon already requires $m_{3/2}$ to be larger than $\sim 10^{-6}\,\mathrm{eV}$ [15, 16], we are forced to conclude that

$$\Lambda \geq 300\,\mathrm{GeV} \tag{26}$$

or, equivalently,

$$m_{3/2} \geq 2.2 \times 10^{-6}\,\mathrm{eV}. \tag{27}$$

In conclusion, we have carefully derived the bounds on the superlight gravitino mass (i.e. the SUSY scale Λ) that follow from SN physics. These limits are completely general in the sense that they do not rely on other particles in a given particular model being light.

Should other particles such as the scalar partners of the goldstino also be light, then the resulting bounds are necessarily tighter. In such clearly less general frame, constraints have also been derived in the literature. Before reviewing these bounds, let us refer to the statement by Brignole *et al* [17] that the amplitude for $e^+e^- \to \tilde{G}\tilde{G}$ based on equation (5) is inconsistent and to the explicit analysis by Clark *et al* [18] showing that the dimension 6 operator $\gamma\tilde{G}\tilde{G}$ derived by Luty and Ponton is actually not present in the effective lagrangian. If true, the bounds derived above are 1 to 2 orders of magnitude worse [19].

Suppose now that the gravitino and S and/or P are very light ($\ll T$). Because now these scalars can be emitted by astrophysical bodies, one has to consider additional energy-loss channels [7, 9].The relevant interaction is given by

$$e^{-1} L_{\mathrm{int}} = -\frac{\kappa}{4}\sqrt{\frac{2}{3}}\left(\frac{m_{\tilde{\gamma}}}{m_{3/2}}\right)(S F^{\mu\nu} F_{\mu\nu} + P \tilde{F}_{\mu\nu} F^{\mu\nu}), \tag{28}$$

where $\kappa \equiv (8\pi)^{1/2} M_{\mathrm{Pl}}^{-1}$.

The main energy drain mechanism is the Primakoff process $\gamma e \to eS/P$ via one photon exchange in the t channel. Calculation of luminosities in the supernova goes along the same lines as before using, of course, the Primakoff cross section and which was derived in reference [7]. The restriction

$$L \leq 10^{52}\,\mathrm{erg} \tag{29}$$

then implies

$$m_{3/2} \geq 30\left(\frac{m_{\tilde{\gamma}}}{100\,\mathrm{GeV}}\right)\mathrm{eV} \tag{30}$$

Since

$$\Gamma(S/P \to 2\gamma) \simeq \frac{\kappa^2}{96\pi}\left(\frac{m_{\tilde{\gamma}}}{m_{3/2}}\right)^2 m_{S/P}^3 \tag{31}$$

sufficiently heavy S/P will decay inside the core. For this NOT to happen, it is required that

$$m_{S/P} \leq 10\,\mathrm{MeV} \tag{32}$$

Also, S/P leave the supernova core without further rescattering provided

$$m_{3/2} \geq 0.3\,\mathrm{eV} \tag{33}$$

(because their mfp is $\sim 10^{10}(m_{3/2}/30\,\mathrm{eV})^2\,\mathrm{cm}$).

For masses below about 0.3 eV, S/P get trapped and their energy is radiated from a S/P-sphere following the Stefan–Boltzmann law $L \propto R^2_{S/P} T^4_{S/P}$. The luminosity L is compatible again with observation (SN1987A) for $m_{3/2} \leq 10^{-1.5}$ eV.

2.2.2 *Limits from the sun* [9]: Should $m_{S/P} < 1$ keV then these particles could be emitted from the sun and the lower limits to their masses are in this case slightly different. In the sun, the mfp of S/P is $\lambda_{S/P} \sim 10^{41}(m_{3/2}/m_{\tilde{\gamma}})^2$ cm and $\lambda_{S/P}$ exceeds the solar radius for $m_{3/2} \geq 10^{-3.5}$ eV (for $m_{\tilde{\gamma}} = 100\,\mathrm{GeV}$).

The emissivity via Primakoff scattering turns out to be

$$\kappa^2 \alpha \left(\frac{m_{\tilde{\gamma}}}{m_{3/2}}\right)^2 T^7 F(\kappa_{\mathrm{D}}/T) V_{\mathrm{Sun}} \tag{34}$$

where the function $F(x)$ (see [1], p. 169) takes care of plasma effects characterized by the Debye momentum κ_{D}. This emissivity is bounded above by the sun luminosity, i.e. 10^{33} erg/s. As a consequence, the gravitino mass should verify

$$m_{3/2} \geq 50\,\mathrm{eV} \tag{35}$$

On the other hand, in the trapping regime (i.e. $m_{3/2} < 10^{-3.5}$ eV), S/P emission is allowed as long as $m_{3/2} \leq 10^{-6}$ eV for in this case their thermal radiation is sufficiently slowed down. Since $(g-2)_\mu$ implies already $m_{3/2} \geq 10^{-6}$ eV [16] we conclude

$$m_{3/2} \geq 50\,\mathrm{eV} \tag{36}$$

for $m_{3/2} < 1$ keV and for $m_{\tilde{\gamma}} \sim O(100)$ GeV.

2.3 *Big bang nucleosynthesis (BBN)*

Primordial Helium-4 abundance depends on the effective degrees of freedom (dof) at nucleosynthesis because $H \sim g_*^{1/2} M_{\mathrm{Pl}}^{-1} T^2$ and, the larger H the sooner do weak interactions decouple. Hence, the neutron to proton ratio is larger and the helium yield rises accordingly.

The effective number of dof for a given species i is

$$g_*^i \binom{\text{fermion}}{\text{boson}} = \binom{1}{7/8} g_i \left(\frac{T_i}{T}\right)^4 \tag{37}$$

Bounds from Y_P are usually presented as bounds on ΔN_ν (i.e. number of extra equivalent neutrino species). It is explicitly given by the formula

$$\Delta N_\nu \binom{\text{fermion}}{\text{boson}} = \binom{8/7}{1} \frac{g_i}{2} \left[\frac{g_*(T_\nu)}{g_*(T_{D_i})}\right]^{4/3}, \tag{38}$$

where T_{D_i} is the decoupling temperature of species i.

The latest analysis from the group in Chicago [20] sets the limit

$$\Delta N_\nu \leq 1 \tag{39}$$

which implies NO bound for the superlight gravitino (equivalent to one neutrino species). That is, the gravitino does not have to decouple from the cosmological thermal bath prior to the BBN era. On the other hand, if the particles S/P are light ($m_{S/P} < 1\,\mathrm{MeV}$, i.e. relativistic at nucleosynthesis) then, they should freeze out before the Universe cools down to $T \sim 200\,\mathrm{MeV}$ so that ΔN_ν is very small (i.e. $g_*(T_\nu)/g_*(T_{Di}) \ll 1$).

S/P are kept in equilibrium through $\gamma e^- \leftrightarrow S/Pe^-$ and this rate falls below the Hubble rate H when

$$\kappa^2 \alpha \left(\frac{m_{\tilde{\gamma}}}{m_{3/2}}\right)^2 T^3 \sim g_*^{1/2} \frac{T^2}{M_{\mathrm{Pl}}}. \tag{40}$$

Taking $T \sim 200\,\mathrm{MeV}$ (and $m_{\tilde{\gamma}} = 100\,\mathrm{GeV}$) we get [9]

$$m_{3/2} \geq 1\,\mathrm{eV}. \tag{41}$$

3. Summary and conclusion

In a wide class of supergravity models with a supersymmetry breaking scale in the TeV range, the gravitino can be very light. In fact, its mass could lie anywhere between $1\,\mu\mathrm{eV}$ and 1 keV. It is also a generic feature of some of the recently considered models that the superlight gravitino is accompanied by a superlight scalar S and pseudoscalar P particles, which are the neutral scalar partners of the goldstino. For energy scales such that $E \gg m_{3/2}$ the longitudinal component of the gravitino dominates and the gravitino effectively behaves as a spin-1/2 Goldstino. The neutral scalars and the gravitino are coupled to matter with strength inversely proportional to the gravitino mass and, hence, they can be abundantly produced in the interior of stellar cores. By requiring that the radiated power does not overcome 10^{52} erg/s in a supernova explosion and 10^{33} erg/s in the case of solar emission, we obtain the following results.

$$m_{3/2} \geq 10^{-5}\,\mathrm{eV} \tag{42}$$

independently of $m_{3/2}$. At this point we should stress that there are by now strong model-independent limits on the gravitino mass ($m > 1.4 \times 10^{-5}\,\mathrm{eV}$) from accelerator data [21] that do not rely on the controversial dimension 6 operator of Luty and Ponton.

Also, in a less general set up,

$$m_{3/2} \geq 30\,\mathrm{eV} \tag{43}$$

or

$$m_{3/2} \leq 0.03\,\mathrm{eV} \tag{44}$$

if $m_{S/P} \leq 10\,\mathrm{MeV}$ and,

$$m_{3/2} \geq 50\,\mathrm{eV} \tag{45}$$

for $m_{S/P} \leq 1\,\mathrm{keV}$.

Finally, from big bang nucleosynthesis arguments we infer the limit

$$m_{3/2} \geq 1\,\text{eV} \tag{46}$$

should $m_{S/P}$ be lighter than about 1 MeV. Results in eqs (43) to (46) are also independent of the coupling in eq. (5).

Acknowledgement

The author is grateful to the organizers of the Workshop for the friendly hospitality and for the nice atmosphere provided in Pune. Work partially supported by the CICYT Research Project AEN95-0882.

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