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DARK MATTER CONSTRAINTS IN HETEROTIC M-THEORY WITH FIVE-BRANE DOMINANCE

D.BAILIN♣ [1], G. V. KRANIOTIS♠ [2] and A. LOVE♠
♣ *Centre for Theoretical Physics,*
University of Sussex,
Brighton BN1 9QJ, U.K.
♠ *Centre for Particle Physics ,*
Royal Holloway and Bedford New College,
University of London,Egham,
Surrey TW20-0EX, U.K.

ABSTRACT

The phenomenological implications of the M-theory limit in which supersymmetry is broken by the auxiliary fields of five-brane moduli is investigated. Assuming that the lightest neutralino provides the dark matter in the universe, constraints on the sparticle spectrum are obtained. Direct detection rates for dark matter are estimated.

[1]D.Bailin@sussex.ac.uk
[2]G.Kraniotis@rhbnc.ac.uk

One of the most interesting developments in M-theory [1, 2, 3, 4, 5, 6, 7, 8, 9, 10] model building is that new non-perturbative tools have been developed which allow the construction of realistic three generation models [11]. In particular the inclusion of five-brane moduli Z_n, (which do not have a weakly coupled string theory counterpart) besides the metric moduli T, S in the effective action leads to new types of $E_8 \times E_8$ symmetry breaking patterns as well as to novel gauge and Kähler threshold corrections. As a result the soft-supersymmetry breaking terms differ substantially from the weakly coupled string.

The phenomenological implications of the effective action of M-theory with the standard embedding of the spin connection into the gauge fields have been investigated in [12, 13, 27, 15, 16]. Some phenomenological implications of non-standard embeddings in M-theory with and without five-branes have been studied in [17, 18].

In a previous letter, we investigated the supersymmetric particle spectrum in the interesting case when the auxiliary fields associated with the five-branes dominate those associated with metric moduli ($F^{Z_n} \gg F^S, F^T$), including the constraint of radiative electroweak symmetry-breaking [19]. It is the purpose of this letter, to extend the calculation and take into account cosmological constraints on the relic abundance of the neutralino assuming it provides the dark matter of the universe in the region of the parameter space in which it is the lightest supersymmetric particle. As we shall see in what follows these constraints are quite restrictive.

The soft supersymmetry-breaking terms are determined by the following functions of the effective supergravity theory [11, 18]:

$$\begin{aligned} K &= -\ln(S+\bar{S}) - 3\ln(T+\bar{T}) + K_5 + \frac{3}{T+\bar{T}}(1+\frac{1}{3}e_O)H_{pq}C_O^p\bar{C}_O^q, \\ f_O &= S + B_O T, \quad f_H = S + B_H T, \\ W_O &= d_{pqr}C_O^p C_O^q C_O^r \end{aligned} \tag{1}$$

where K is the Kähler potential, W_O the observable sector perturbative superpotential, C_O^p are observable sector matter fields and f_O, f_H are the gauge kinetic functions for the observable and hidden sector gauge groups respectively. K_5 is the Kähler potential for the five-brane moduli Z_n and H_{pq} is some T-independent metric. Also

$$e_O = b_O\frac{T+\bar{T}}{S+\bar{S}}, \quad e_H = b_H\frac{T+\bar{T}}{S+\bar{S}} \tag{2}$$

and the coefficients $b_{O,H}, B_{O,H}$ are given in terms of the instanton numbers $\beta_{O,H}$ and the five brane charges β_n by the following expressions

$$\begin{aligned} b_O &= \beta_O + \sum_{n=1}^{N}(1-z_n)^2\beta_n \\ B_O &= \beta_O + \sum_{n=1}^{N}(1-Z_n)^2\beta_n \\ B_H &= \beta_H + \sum_{n=1}^{N}(Z_n)^2\beta_n \end{aligned}$$

$$b_H = \beta_H + \sum_{n=1}^{N} z_n^2 \beta_n \tag{3}$$

and the five-brane moduli are denoted by Z_n whose $\mathrm{Re}Z_n \equiv z_n = \frac{x_n}{\pi\rho} \in (0,1)$ are the five-brane positions in the normalized orbifold coordinates. Since a Calabi-Yau manifold is compact, the net magnetic charge due to orbifold planes and 5-branes is zero. Consequently the following cohomology condition is satisfied

$$\beta_O + \sum_{n=1}^{N} \beta_n + \beta_H = 0 \tag{4}$$

S, T are the dilaton and Calabi-Yau moduli fields and C^p charged matter fields. The superpotential and the gauge kinetic functions are exact up to non-perturbative effects.

Given eqs(1) one can determine [11, 18] the soft supersymmetry breaking terms for the observable sector gaugino masses $M_{1/2}$, scalar masses m_0 and trilinear scalar couplings A as functions of the auxiliary fields F^S,F^T, F^n of the moduli S, T fields and five-brane moduli Z_n respectively.

$$\begin{aligned}
M_{1/2} &= \frac{1}{(S+\bar{S})(1+\frac{B_O T + \bar{B}_O \bar{T}}{S+\bar{S}})}(F^S + F^T B_O + T F^n \partial_n B_O) \\
m_0^2 &= V_0 + m_{3/2}^2 - \frac{1}{(3+e_O)^2}\Big[e_O(6+e_O)\frac{|F^S|^2}{(S+\bar{S})^2} \\
&+ 3(3+2e_O)\frac{|F^T|^2}{(T+\bar{T})^2} - \frac{6e_O}{(S+\bar{S})(T+\bar{T})} Re F^S \bar{F}^{\bar{T}} \\
&+ \Big(\frac{e_O}{b_O}(3+e_O)\partial_n\partial_{\bar{m}} b_O - \frac{e_O^2}{b_O^2}\partial_n b_O \partial_{\bar{m}} b_O\Big) F^n \bar{F}^{\bar{m}} \\
&- \frac{6e_O}{b_O}\frac{\partial_{\bar{n}} b_O}{S+\bar{S}} Re F^S \bar{F}^{\bar{n}} + \frac{6e_O}{b_O}\frac{\partial_{\bar{n}} b_O}{T+\bar{T}} Re F^T \bar{F}^{\bar{n}}\Big], \\
A &= -\frac{1}{3+e_O}\Big\{\frac{F^S(3-2e_O)}{S+\bar{S}} + \frac{3e_O F^T}{T+\bar{T}} \\
&+ F^n\Big(3\frac{e_O}{b_O}\partial_n b_O - (3+e_O)\partial_n K_5\Big)\Big\}
\end{aligned} \tag{5}$$

where $\partial_n \equiv \frac{\partial}{\partial Z_n}$. The bilinear B-parameter associated with a non-perturbatively generated μ term in the superpotential is given by [18]:

$$\begin{aligned}
B_\mu &= \frac{F^S(e_O - 3)}{(3+e_O)(S+\bar{S})} - \frac{3(e_O+1)F^T}{(T+\bar{T})(3+e_O)} \\
&+ \frac{1}{3+e_O}\Big[(3+e_O)F^n\partial_n K_5 - 2F^n\frac{e_O}{b_O}\partial_n b_O\Big] - m_{3/2}
\end{aligned} \tag{6}$$

From now on we assume that only one five-brane contributes to supersymmetry-breaking [3]. Then the auxiliary fields are given by [18, 22]

$$\begin{aligned} F^1 &= \sqrt{3} m_{3/2} C(\partial_1 \partial_{\bar{1}} K_5)^{-1/2} \sin\theta_1 \\ F^S &= \sqrt{3} m_{3/2} C(S+\bar{S}) \sin\theta \cos\theta_1 \\ F^T &= m_{3/2} C(T+\bar{T}) \cos\theta \cos\theta_1 \end{aligned} \tag{7}$$

The goldstino angles are denoted by θ, θ_1, $m_{3/2}$ is the gravitino mass and $C^2 = 1 + \frac{V_0}{3m_{3/2}^2}$ with V_0 the tree level vacuum energy density. The five-brane dominated supersymmetry-breaking scenario corresponds to $\theta_1 = \frac{\pi}{2}$, i.e $F^T, F^S = 0$, and we take the five brane which contributes to supersymmetry breaking to be located at $z_1 = 1/2$ in the orbifold interval. We also set $C = 1$ in the above expressions assuming zero cosmological constant.

The resulting supersymmetric particle spectrum for a single five-brane present with $z_1 = \frac{1}{2}$ has been investigated in [19]. Our parameters are, e_O, $\partial_1 \partial_{\bar{1}} K_5$, $\partial_1 K_5$, $m_{3/2}$, $sign\ \mu$ (which is not determined by the radiative electroweak symmetry breaking constraint), where μ is the Higgs mixing parameter in the low energy superpotential. The ratio of the two Higgs vacuum expectation values $\tan\beta = \frac{<H_2^0>}{<H_1^0>}$ is also a free parameter if we leave B to be determined by the minimization of the one-loop Higgs effective potential. If B instead is given by (6), one determines the value of $\tan\beta$. For this purpose we take μ independent of T and S because of our lack of knowledge of μ in M-theory. We treat e_O as a free parameter as the problem of stabilizing the dilaton and other moduli has not yet been solved, although there has been an interesting work in this area [23].

The instanton numbers are model dependent. In this paper we choose to work with the interesting example [18] with $\beta_O = -2$ and $\beta_1 = 1$ which implies $b_O = -7/4$. This implies that $b_H = 5/4$ and allow us to study the region of parameter space with $-1 < e_O \le 0$, which is not accessible in strongly coupled M-theory scenarios with standard embedding. We also choose $\partial_1 K_5 = \partial_1 \partial_{\bar{1}} K_5 = 1$. In an earlier paper [19] we have investigated deviations from these values and have verified the robustness of our results.

We use the following experimental bounds from unsuccessful searches at LEP and Tevatron for supersymmetric particles [20]. We require the lightest chargino $M_{\chi_1^+} \geq 90$ GeV, and the lightest Higgs, $m_{h_0} \geq 83$ GeV. A lower limit on the mass of the lightest stop $m_{\tilde{t}_2} > 86$ GeV, from $\tilde{t}_2 \to c\chi_1^0$ decay in D0 is imposed. The stau mass eigenstate $(\tilde{\tau})$ should be heavier than 81 GeV from LEP2 results.

The soft masses start running from a mass $R_{11}^{-1} \sim 7.5 \times 10^{15}$ GeV with R_{11} the extra M-theory dimension. Then using (5),(6) as boundary conditions for the soft terms, one evolves the renormalization group equations down to the weak scale and determines the sparticle spectrum compatible with the constraints of correct electroweak symmetry breaking and the above experimental constraints on the sparticle spectrum.

[3] We assume negligible CP-violating phases in the soft terms.

Electroweak symmetry breaking is characterized by the extrema equations

$$\begin{aligned}\frac{1}{2}M_Z^2 &= \frac{\bar{m}_{H_1}^2 - \bar{m}_{H_2}^2 \tan^2\beta}{\tan^2\beta - 1} - \mu^2 \\ -B\mu &= \frac{1}{2}(\bar{m}_{H_1}^2 + \bar{m}_{H_2}^2 + 2\mu^2)\sin 2\beta \end{aligned} \tag{8}$$

where

$$\bar{m}_{H_1,H_2}^2 \equiv m_{H_1,H_2}^2 + \frac{\partial \Delta V}{\partial v_{1,2}^2} \tag{9}$$

and $\Delta V = (64\pi^2)^{-1}\mathrm{STr}M^4[ln(M^2/Q^2) - \frac{3}{2}]$ is the one loop contribution to the Higgs effective potential. We include contributions only from the third generation of particles and sparticles.

Since $\mu^2 \gg M_Z^2$ for most of the allowed region of the parameter space [24], the following approximate relationships hold at the electroweak scale for the masses of neutralinos and charginos, which of course depend on the details of electroweak symmetry breaking.

$$\begin{aligned} m_{\chi_1^\pm} \sim m_{\chi_2^0} \sim 2m_{\chi_1^0} \\ m_{\chi_{3,4}^0} \sim m_{\chi_2^\pm} \sim |\mu| \end{aligned} \tag{10}$$

In (10) $m_{\chi_{1,2}^\pm}$ are the chargino mass eigenstates and $m_{\chi_i^0}, i = 1\ldots 4$ are the four neutralino mass eigenstates with $i = 1$ denoting the lightest neutralino. The former arise after diagonalization of the mass matrix.

$$M_{ch} = \begin{pmatrix} M_2 & \sqrt{2}m_W \sin\beta \\ m_W \cos\beta & -\mu \end{pmatrix} \tag{11}$$

where M_2 denotes the weak gaugino mass and M_1 will denote the $U(1)_Y$ gaugino (Bino) mass. The stau mass matrix is given by the expression

$$\mathcal{M}_\tau^2 = \begin{pmatrix} \mathcal{M}_{11}^2 & m_{\tilde{\tau}}(A_\tau + \mu\tan\beta) \\ m_{\tilde{\tau}}(A_\tau + \mu\tan\beta) & \mathcal{M}_{22}^2 \end{pmatrix} \tag{12}$$

where $\mathcal{M}_{11} = m_L^2 + m_\tau^2 - \frac{1}{2}(2M_W^2 - M_Z^2)\cos 2\beta$ and $\mathcal{M}_{22} = m_E^2 + m_\tau^2 + (M_W^2 - M_Z^2)\cos 2\beta$. where m_L^2, m_E^2 refer to scalar soft masses for lepton doublet, singlet respectively.

As has been first noted in [18], in the case when only the five-branes contribute to supersymmetry-breaking the ratio of scalar masses to gaugino mass, $m_0/|M_{1/2}| > 1$ for $e_O > -0.65$. This is quite interesting since scalar masses larger than gaugino masses are not easy to obtain in the weakly-coupled heterotic string or M-theory compactification with standard embedding. As we shall see cosmological constraints become important in this region of the parameter space. In the case of non-standard embeddings without five-branes there is small region of the parameter space where is possible to have $m_0/|M_{1/2}| > 1$.

Assuming R-parity conservation the LSP is stable, and consequently if it is neutral can provide a good dark matter candidate. We assume that the dark matter is in the form of neutralinos. The lightest neutralino is a linear combination of the superpartners of the photon, Z^0 and neutral-Higgs bosons,

$$\chi_1^0 = N_{11}\tilde{B} + N_{12}\tilde{W}^3 + N_{13}\tilde{H}_1^0 + N_{14}\tilde{H}_2^0 \tag{13}$$

The neutralino 4×4 mass matrix can be written as

$$\begin{pmatrix} M_1 & 0 & -M_Z A_{11} & M_Z A_{21} \\ 0 & M_2 & M_Z A_{12} & -M_Z A_{22} \\ -M_Z A_{11} & M_Z A_{12} & 0 & \mu \\ M_Z A_{21} & -M_Z A_{22} & \mu & 0 \end{pmatrix}$$

with

$$\begin{pmatrix} A_{11} & A_{12} \\ A_{21} & A_{22} \end{pmatrix} = \begin{pmatrix} \sin\theta_W \cos\beta & \cos\theta_W \cos\beta \\ \sin\theta_W \sin\beta & \cos\theta_W \sin\beta \end{pmatrix}$$

When the observational data on temperature fluctuations, type Ia supernovae, and gravitational lensing are combined with popular cosmological models, the dark matter relic abundance (Ω_{LSP}) typically satisfies [26]

$$0.1 \le \Omega_{LSP} h^2 \le 0.4 \tag{14}$$

where h is the reduced Hubble constant.

We calculated the relic abundance of the lightest neutralino in the scenarios we have considered using standard techniques [25]. When these results are confronted with the (model-dependent) bounds (14) derived from the observational data further constraints on the parameters $m_{3/2}, \tan\beta, \mu, e_O$ are obtained and these give new constraints on the sparticle spectrum.

In figs.(1,2,3) we display the relic abundance of the lightest neutralino versus $\tan\beta$ for different values of the gravitino mass and the parameter e_O. We see that $\Omega_{LSP} h^2 \le 0.4$ puts $m_{3/2}$ dependent lower bounds on the values of $\tan\beta$. Let us start the discussion with the case $e_0 = -0.6, \mu < 0$, fig.(1). In this case the upper limit on the relic abundance provides the following lower bounds on $\tan\beta$. For instance, for $m_{3/2} = 170$GeV $\tan\beta > 5$ while for $m_{3/2} = 230$GeV $\tan\beta > 20$. Values of the gravitino mass $m_{3/2} < 170$GeV come into contradiction with the lower experimental bounds on the lightest Higgs mass and lightest chargino mass imposed from unsuccessful searches for supersymmetric particles in accelerator experiments. The lower limit on the relic abundance ($\Omega_{LSP} h^2 > 0.1$) imposes further constraints on the gravitino mass for $\tan\beta \ge 26$. In particularfor $\tan\beta = 26$, $m_{3/2} \ge 182GeV$ which results in: $m_{\chi_1^+} \ge 102GeV, m_{h^0} \ge 115GeV$. Similarly for $\tan\beta = 28$, $m_{3/2} \ge 210$GeV and

$\tan\beta$	$m_{3/2}^{max}$	m_{h^0}	$m_{\chi_1}^{+}$	$m_{\chi_1^0}$	$m_{\tilde{t}_2}$	$m_{\tilde{\tau}_2}$
22	380GeV	118.5 GeV	99 GeV	54 GeV	155 GeV	310 GeV
28	381GeV	118 GeV	99 GeV	54 GeV	154 GeV	284 GeV
30	400GeV	119 GeV	104 GeV	56 GeV	164 GeV	288 GeV
32	495GeV	122 GeV	131 GeV	70 GeV	216 GeV	347 GeV
34	750GeV	129 GeV	201 GeV	106 GeV	371 GeV	513. GeV

Table 1: Upper bounds on sparticle masses resulting from Eq. (14) for $e_O = -0.4, \mu < 0$ for various values of $\tan\beta$.

$m_{\chi_1^+} \geq 119.5$GeV. In this case the lightest Higgs mass $m_{h^0} \geq 117.6$GeV. The allowed values of $m_{3/2}$ as a function of $\tan\beta$, compatible with (14), and for $e_O = -0.6, \mu < 0$ are shown in fig.4. Also shown in Figs.5,4 are upper bounds on the lightest chargino, lightest Higgs and lightest neutralino masses respectively, compatible with the upper cosmological limit on the relic abundance for both signs of μ. The resulting model can be tested at accelerator experiments.

The cosmological constraints become even more important when $e_O \to 0$. In figs. 2 and 3 we plot the relic abundance of the lightest neutralino versus $\tan\beta$ for different values of the gravitino mass for $e_O = -0.4$. The lower experimental bounds (from unsuccessful searches in accelerator experiments) on the lightest chargino mass and on the lightest Higgs mass now require that $m_{3/2} \geq 380GeV$. For $\mu < 0$ we see that $\Omega_{LSP}h^2 \leq 0.4$ requires $\tan\beta \geq 22$ for any allowed value of the gravitino mass and for $22 \leq \tan\beta \leq 30$, the relic abundance of the lightest neutralino is in the range:$0.05 \leq \Omega_{LSP}h^2 \leq 0.33$. The upper bounds on the gravitino, lightest chargino, lightest Higgs, lightest stop and lightest stau masses as a function of $\tan\beta$ compatible with the upper limit on the relic abundance are summarised in Table 1.

In fig. (6) we plot the relic abundance of the lightest neutralino versus e_O for $\tan\beta = 15, 26$ respectively and fixed Bino mass at the unification scale $M_1(M_U) = 126$GeV. This choice corresponds to a lightest neutralino mass of ~ 54GeV corresponding to a lightest chargino mass of about 100 GeV, the current experimental lower bound. From fig.6 we observe that for $e_O \geq -0.55$, $\tan\beta = 15$, $\Omega_{LSP}h^2 > 0.4$ and for $e_O > -0.5$ is greater than 1. On the other hand for $e_O \to -1$ the cosmological constraints are more easily satisfied although for e_O too close to -1 $\Omega_{LSP}h^2 > 0.1$ becomes more difficult to satisfy. Also for $e_O < -0.8$ the lightest stau becomes the lightest supersymmetric particle. For $\tan\beta = 26$, we have that for $e_O \geq -0.45$, $\Omega_{LSP}h^2 > 0.4$.

Assuming that the neutralinos provide the cold dark matter in our Galaxy we calculated its direct detection rates for various nuclei. For an LSP moving with velocity v_z with respect to the detector nuclei the detection rate for a target with mass m is given by [25]

$$R = \frac{\rho^{0.3}_{\chi^0_1}}{m_{\chi^0_1}} \frac{m}{Am_p} \int f(\mathbf{v})|\mathbf{v_z}|\sigma(|\mathbf{v}|)\mathbf{d^3v}, \tag{15}$$

where $\rho^{0.3}_{\chi^0_1}$ denotes the local LSP mass density normalized to the standard value of $0.3\mathrm{GeVcm}^{-3}$, $f(\mathbf{v})$ a Maxwell velocity distribution and σ denotes the neutralino-nucleous elastic cross section.

For $e_0 = -0.6$ the cross sections of neutralinos with nucleon result in large total detection rates in the cosmologically interesting region for large values of $\tan\beta$. In particular, for $^{73}Ge,^{208}Pb,^{131}Xe$ detectors, detection rates of the neutralinos are in the range of order $10^{-3} - O(1)$events/Kg/day for $\mu < 0$. The larger total event rates occur for $\tan\beta \geq 24$. The results for $\mu > 0$ are similar. This illustrates the fact that $\Omega_{LSP}h^2 \sim \frac{10^{-37}cm^2}{<\sigma_{anni}v>}$ and the neutralino annihilation cross section is roughly proportional to the neutralino scattering cross section. Thus as the LSP abundance decreases, its scattering cross section generally increases. For $\Omega_{LSP}h^2 \sim 0.1$ this results in an increased event rate. For values of $e_O > -0.6$ the total detection rates are smaller. This behaviour of the detection rates can be understood by investigating the neutralino-nucleon scalar (spin-independent) cross section, which in this model is the dominant contribution to the total neutralino-nucleous elastic cross section.

The scalar nucleon-LSP cross section is given by [12, 29]

$$\sigma^{(nucleon)}_{scalar} = \frac{8G_F^2}{\pi} M_W^2 m_{red}^2 \left[\frac{G_1(h_0) I_{h_0}}{m_{h_0}^2} + \frac{G_2(H) I_H}{m_H^2} + \cdots \right]^2 \tag{16}$$

where

$$\begin{aligned} G_1(h_0) &= (-N_{11}\tan\theta_W + N_{21})(N_{31}\sin\alpha + N_{41}\cos\alpha) \\ G_2(H) &= (-N_{11}\tan\theta_W + N_{21})(-N_{31}\cos\alpha + N_{41}\sin\alpha) \end{aligned} \tag{17}$$

and

$$I_{h_0,H} = \sum_q l_q^{h_0,H} m_q < N|\bar{q}q|N > \tag{18}$$

and

$$\begin{aligned} l_q^{h_0} &= \frac{\cos\alpha}{\sin\beta} \quad l_q^H = \frac{\sin\alpha}{\sin\beta} \quad \text{for} \quad q = u, c, t \\ l_q^{h_0} &= -\frac{\sin\alpha}{\cos\beta} \quad l_q^H = \frac{\cos\alpha}{\cos\beta} \quad \text{for} \quad q = d, s, b \end{aligned} \tag{19}$$

where the two first terms inside the brackets refer to the diagrams with h_0 and H-exchanges in the t-channel and the the ellipsis refers to the graphs with squark-exchanges in the s- and u-channels [12]. In equation (16) m_{red} is the neutralino-nucleon reduced mass, h_0, H denote the lightest Higgs and CP-even heavier Higgs respectively and α is the Higgs mixing angle [4].

[4] We determine the Higgs mixing angle numerically by diagonalizing the one-loop CP-even Higgs mass matrix.

Of particular interest is the $\tan\beta$ dependence of the scalar neutralino-nucleon cross section $\sigma_{scalar}^{nucleon}$. For high values of $\tan\beta$ the corresponding cross section generically increases (see Figs.7, 8). The calculated cross section for high $\tan\beta$ reaches the sensitivity of current dark matter experiments for the nucleon-neutralino scalar cross section in the range 1×10^{-9}nb$(2.6 \times 10^{-15} GeV^{-2}) \leq \sigma_{scalar}^{(nucleon)} \leq 3 \times 10^{-8}nb(7.7 \times 10^{-14} GeV^{-2})$ [29]. The larger the value of the e_O parameter the smaller the cross-section for fixed $\tan\beta$, signμ and fixed bino mass. Also the cross section decreases for increasing values of the gravitino mass for fixed e_O and signμ as is evident from figs.(7, 8). The two upper curves in graphs 7 and 8 are cross-sections for an LSP mass about 54GeV. The corresponding spin-dependent cross sections [12] are much smaller by two to three orders of magnitude.

Recently, large detection rates have been obtained in type I string theories formulated as orientifold compactifications of type IIB string theory [27, 28]. In particular, in the mirage unification scenario large detection rates have been obtained. This scenario, differs from the ones studied in this paper in the following respects:First in the mirage unification scenario the lightest neutralino has a large Higgsino component while in the 5-brane dominated limit it is almost a Bino. Thus in the first case besides a large scalar cross-section the LSP has also a rather large spin-dependent couplings with the nuclei.

Second, in the type I model, the nucleon-neutralino cross section is large and consequently the detection rates are large when $\tan\beta$ is small, i.e. $\tan\beta \leq 8$. Also in this case the high neutralino-nucleon cross sections correspond to relatively low relic neutralino densities, i.e. $\Omega_{LSP} h^2 \leq 0.1$ and therefore another form of dark matter might be needed to close the Universe. As we saw in the current model large cross-sections occur in the high $\tan\beta$ region. Thus the two models lead to different predictions.

As well as predictions for direct detection for the lightest neutralino χ_1^0 we have obtained bounds on the value of $\tan\beta$ from the cosmological constraints on the relic abundance. We have also calculated the maximum lightest chargino mass and the lightest maximum Higgs mass as a function of $\tan\beta$ for various values of the ratio of the scalar masses to gaugino masses. The resulting sparticle spectra should be tested in accelerator experiments in Tevatron and LHC.

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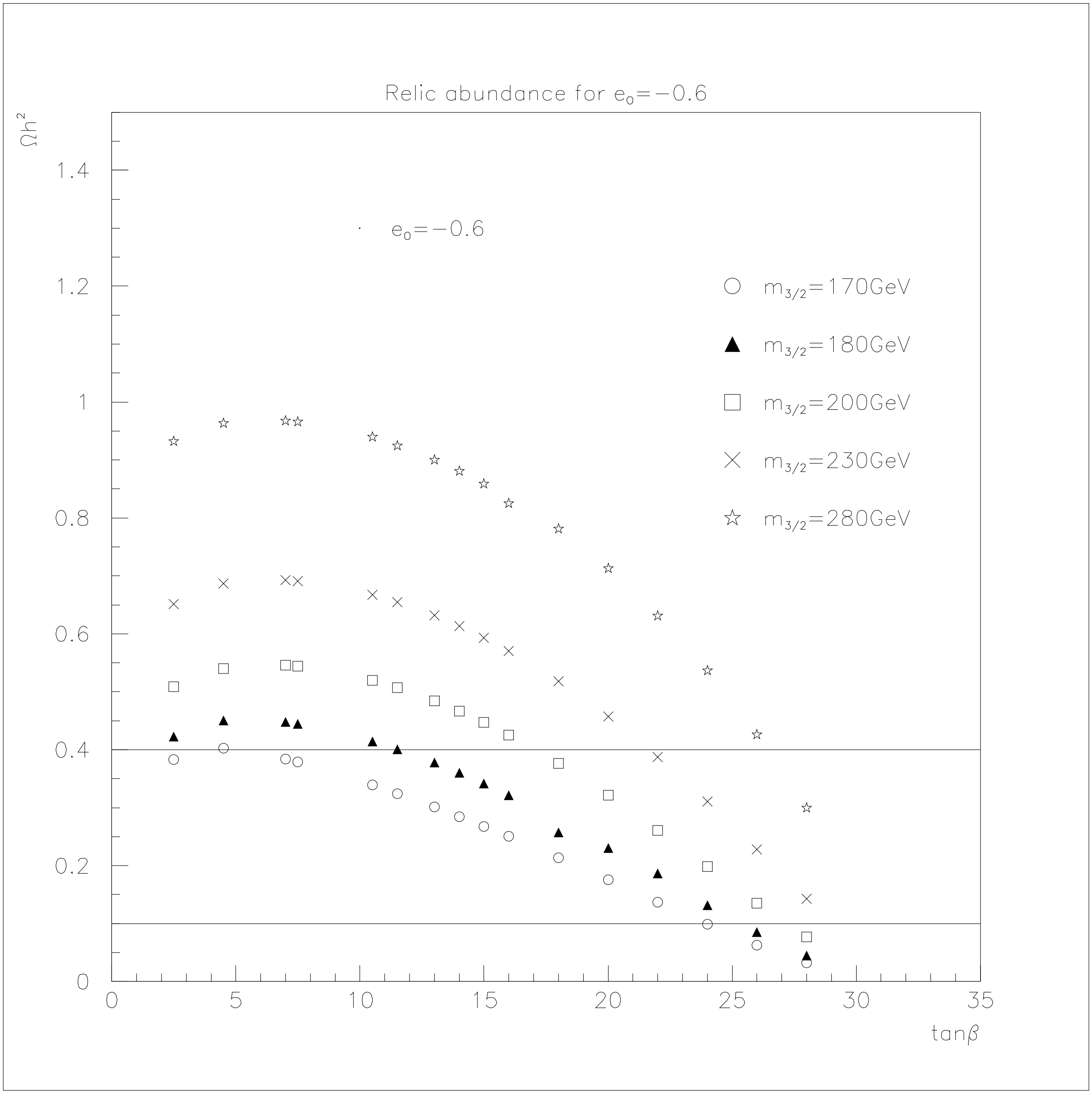

Figure 1: Relic abundance of the LSP versus $\tan\beta$ for various values of the gravitino mass and $e_O = -0.6, \mu < 0$. We also exhibit the upper and lower cosmological bounds on the LSP relic abundance

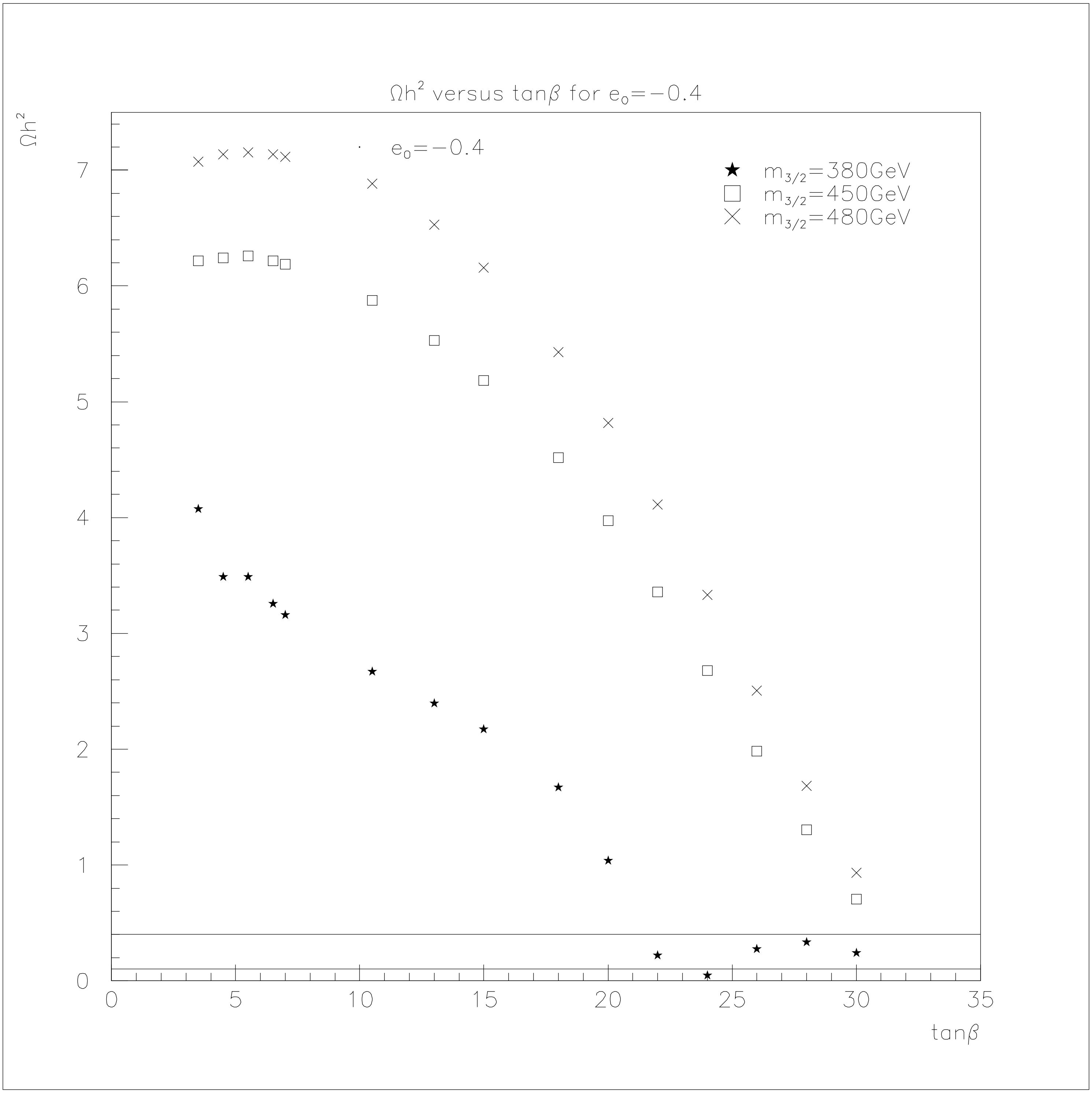

Figure 2: Relic abundance of LSP vs $\tan\beta$ for $e_O = -0.4, \partial_1 K_5 = \partial_1 \partial_{\bar{1}} K_5 = 1, m_{3/2} = 380GeV, 450GeV, 480GeV, \mu < 0$.

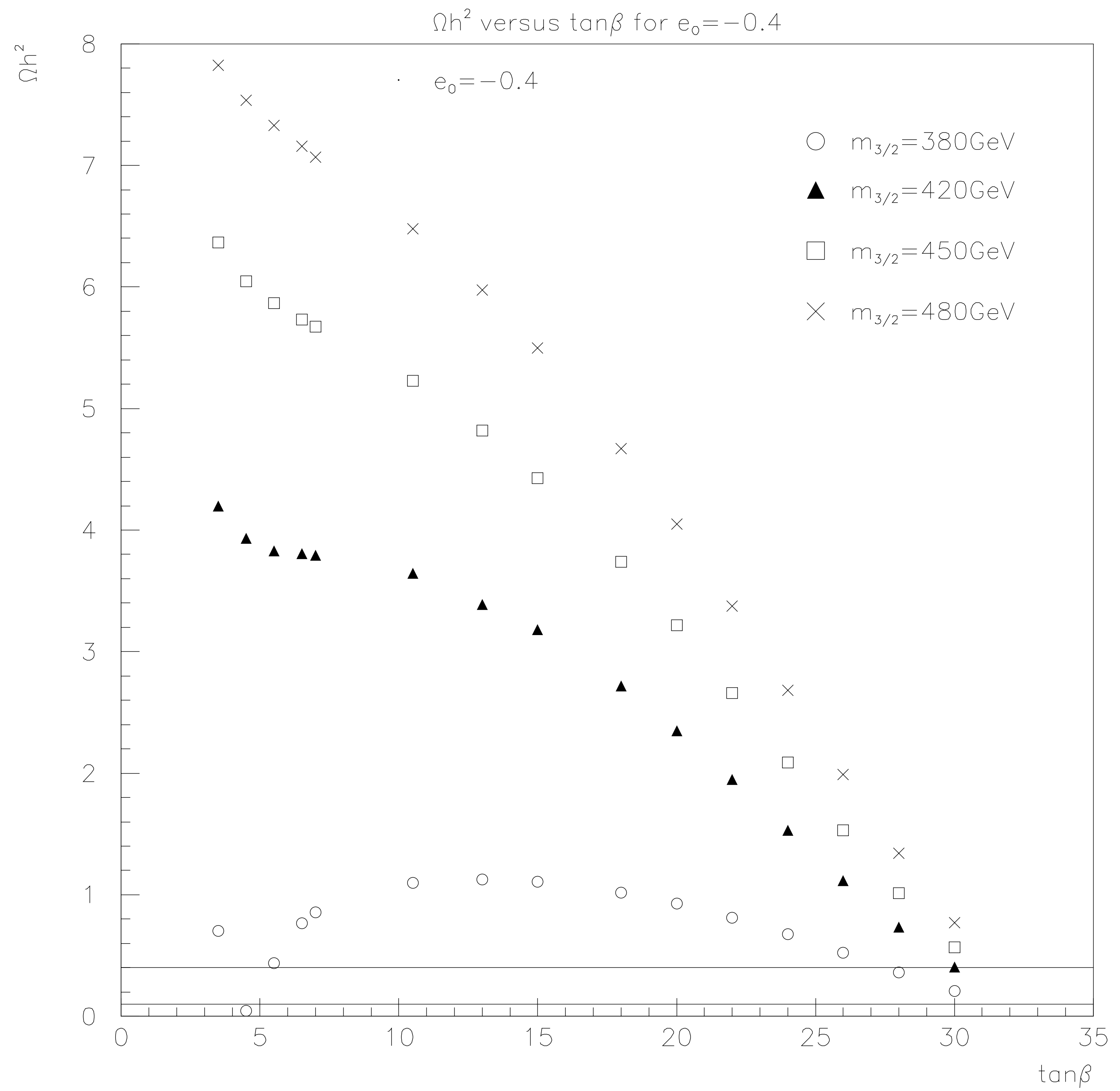

Figure 3: Relic abundance of LSP vs $\tan\beta$ for $e_O = -0.4, \partial_1 K_5 = \partial_1\partial_{\bar{1}} K_5 = 1, m_{3/2} = 380GeV, 420GeV, 450GeV, 480GeV \mu > 0.$

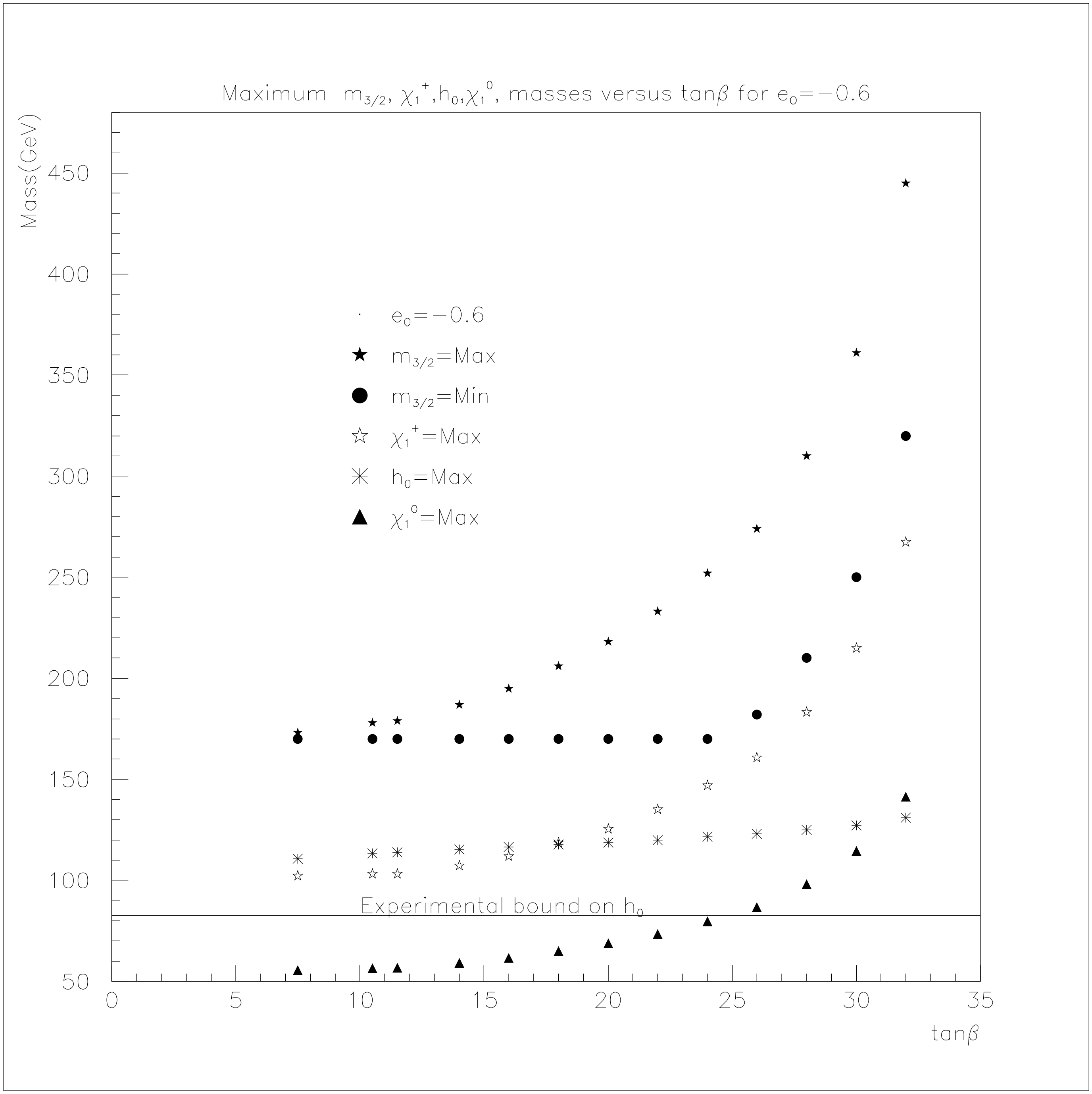

Figure 4: Maximum gravitino, lightest chargino and lightest Higgs masses (imposed by $\Omega h^2 \leq 0.4$) versus $\tan\beta$ for $e_O = -0.6$, $\mu < 0$.

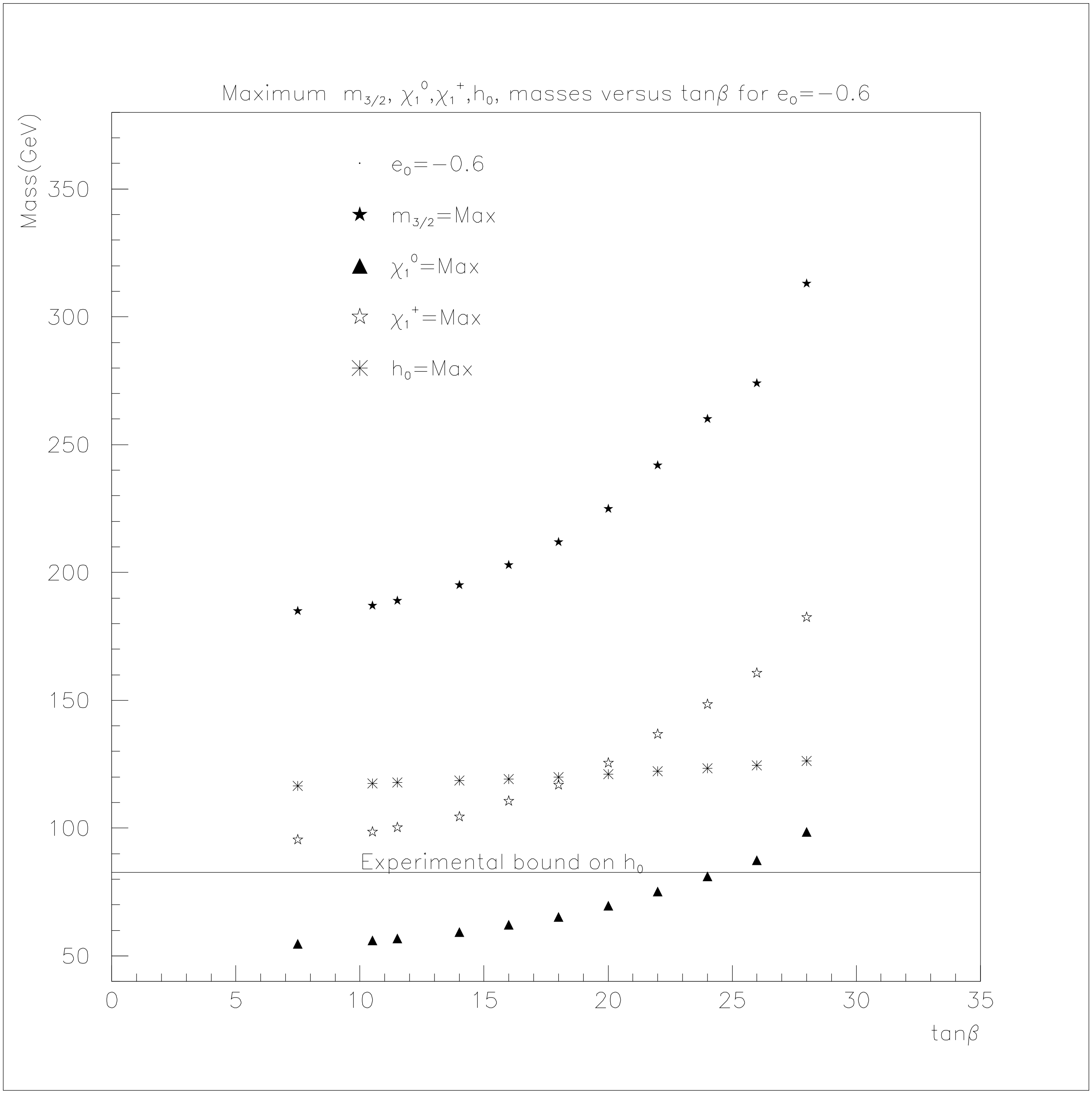

Figure 5: Maximum gravitino, lightest chargino and lightest Higgs masses (imposed by $\Omega h^2 \leq 0.4$) versus $\tan\beta$ for $e_O = -0.6$, $\mu > 0$.

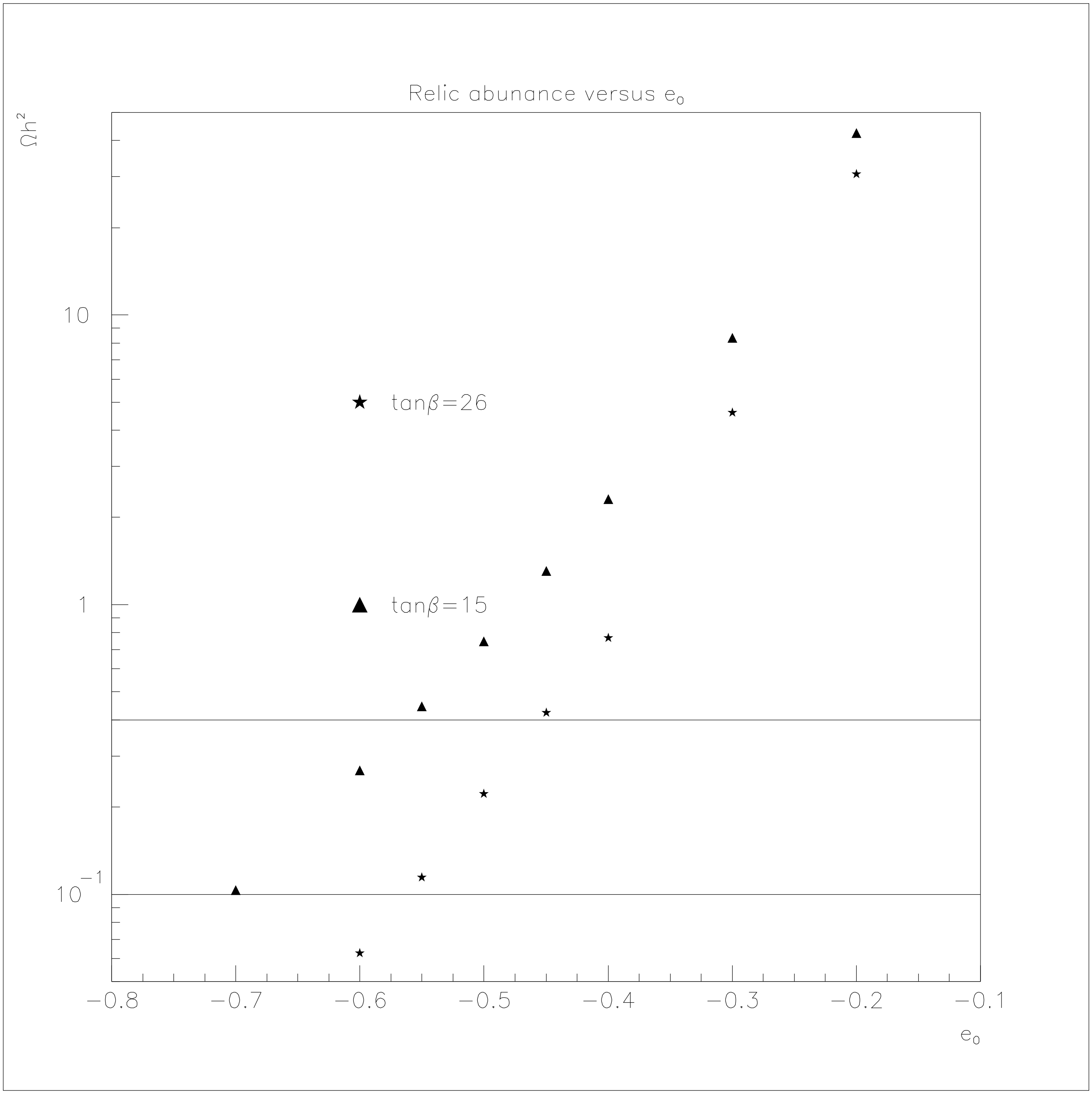

Figure 6: Relic abundance of LSP vs e_O for $\tan\beta = 15, 26, \partial_1 K_5 = \partial_1\partial_{\bar{1}} K_5 = 1, |M_1(M_U)| = 126 GeV, \mu < 0$.

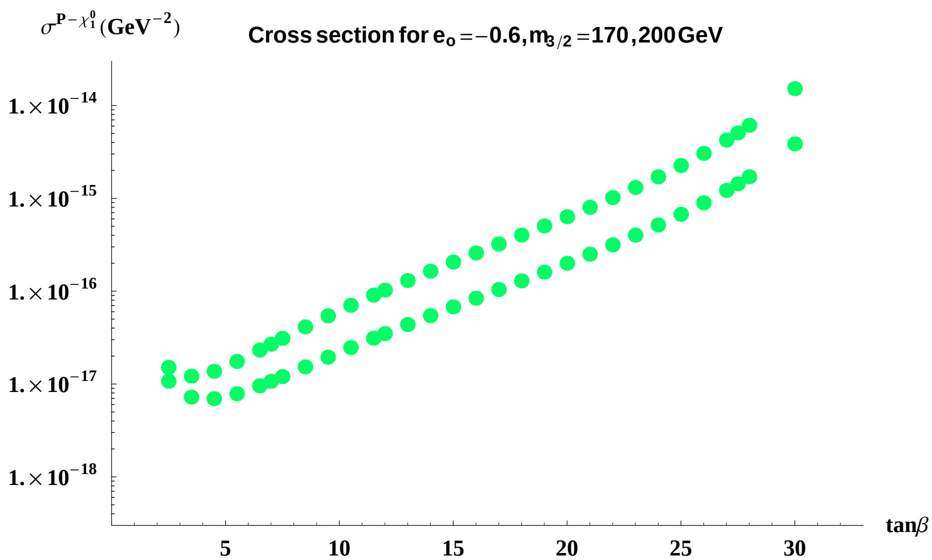

Figure 7: Proton-scalar LSP cross section versus $\tan\beta$ for fixed $e_o = -0.6, \mu < 0, m_{3/2} = 170 (upper curve), 200 GeV (lower curve)$.

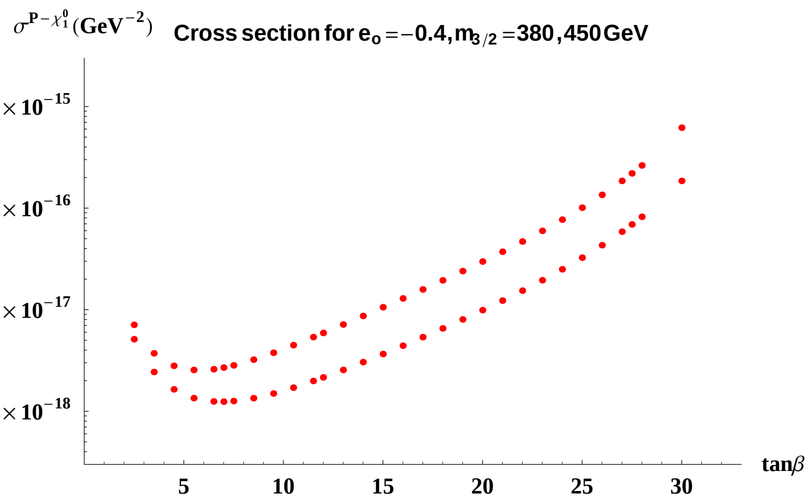

Figure 8: Proton-LSP cross section versus $\tan\beta$ for $e_o = -0.4, \mu < 0, m_{3/2} = 380GeV(upper\ curve), 450GeV(lower\ curve)$.