APCTP 97–21
TIFR/TH/97–62
YUMS–97/30
hep-ph/yymmnn
December 1997

Stringent New Bounds on Supersymmetric Higgs Bosons from Existing Tevatron Data

Manuel Drees[1], Monoranjan Guchait[2,3], and Probir Roy[3]
[1] *APCTP, 207–43 Cheongryangryi–dong, Tongdaemun–gu, Seoul 130–012, Korea*
[2] *Physics Department, Yonsei University, Seoul 120–749, Korea*
[3] *Tata Institute of Fundamental Research, Mumbai, India*

Abstract

Tight limits are derived on models with extended Higgs sectors, including the Minimal Supersymmetric Standard Model (MSSM), by exploiting the data on $\tau^+\tau^- +$ 2 jets final states, used by the CDF collaboration to place limits on third generation leptoquarks. The main observation here is that both leptoquark production and associated $b\bar{b}$ Higgs production can lead to $b\bar{b}\tau^+\tau^-$ final states. For the MSSM we find that CP–odd neutral Higgs boson masses up to 190 GeV are excluded if the ratio of vacuum expectation values tanβ is 50.

Supersymmetry is widely regarded to be the most likely cure of various ills plaguing the Standard Model (SM) [1]. Supersymmetric field theories not only contain superpartners for all known SM particles, they also need (at least) two $SU(2)$ doublets of Higgs superfields to give masses to both up–type and down–type quarks [2]. After electroweak symmetry breaking, this leaves (at least) five physical Higgs bosons: two neutral CP–even scalars h, H, with h being the lighter one; a neutral CP–odd scalar A; and charged Higgs bosons $H^{\pm}$. At the tree level the masses and mixings of these five states are determined by the values of just two parameters [3], often taken to be the mass m_A of the CP–odd Higgs boson and the ratio of vacuum expectation values $\tan\beta \equiv \langle H_u^0 \rangle / \langle H_d^0 \rangle$, where H_u is the Higgs doublet coupling to up–type quarks. The masses and mixings of the two neutral CP–even scalars can receive substantial radiative corrections [4], the dominant contribution coming from loops involving top quarks and their superpartners.

At present, the most stringent limits on the Higgs sector of the MSSM come from the e^+e^- collider LEP at CERN. Unsuccessful searches for Zh and Ah production lead to the mass bounds [5] $m_h, \ m_A \geq 75$ GeV for $\tan\beta > 1$. If $m_A^2 \gg M_Z^2$, the lower bound on m_h increases to 88.6 GeV. If $\tan\beta < 1$ or $\tan\beta > 40$, one can derive additional bounds from the fact that the branching ratio for top quark decay into charged Higgs bosons cannot be large [6], given the good agreement between SM predictions for $\sigma\,(p\bar{p} \rightarrow t\bar{t}X) \cdot \mathrm{B}\,(t \rightarrow Wb \rightarrow l\nu b)$ and the measurements [7] of this quantity. However, the combination of these bounds still leaves a large region of the parameter space unexplored.

The most widely studied mechanism for producing neutral Higgs bosons at the Tevatron is the associated production of a CP–even scalar with a W or Z gauge boson [8]. Since two quite massive particles have to be produced in the final state, the total cross section is small. In addition, one will probably have to require the gauge boson to decay leptonically; this reduces the rate by another factor 0.22 (0.06) for W (Z)+ Higgs production. As a result, these processes will not be useful until an integrated luminosity of several fb^{-1} has been accumulated.

Instead, in this Letter the associated hadroproduction of a neutral Higgs boson with a $b\bar{b}$ pair is addressed in the large $\tan\beta$ regime. Here we make use of the fact that the $Ab\bar{b}$ coupling is proportional to $\tan\beta$, i.e. the cross section grows as $\tan^2\beta$. This can easily lead to a factor 1,000 enhancement over the corresponding cross section [9] for SM Higgs production. In addition, in the region $\tan\beta \gg 1$ one of the two neutral CP–even Higgs bosons of the MSSM is always nearly degenerate with A, and its coupling to $b\bar{b}$ pairs has essentially the same strength as the $Ab\bar{b}$ coupling. For small m_A this near–degenerate state is the light Higgs boson h, while for larger m_A it is H. At the tree level, the cross–over occurs at $m_A = M_Z$. After loop corrections are taken into account, the cross–over occurs at the upper bound on m_h, the precise value of which depends on the details of the sparticle spectrum, in particular on the parameters of the stop mass matrix; for squark masses not exceeding 1 TeV, it is generally between 100 and 125 GeV. However, the identity of this near–degenerate state is of little practical importance; all that matters is that it increases our cross section by essentially a factor of two, and that $b\bar{b}A$ and $b\bar{b}h/H$ events have the same kinematical features.

Our cross section calculation is based on matrix elements for the $2 \rightarrow 3$ subprocesses $gg \rightarrow b\bar{b}$ Higgs and $q\bar{q} \rightarrow b\bar{b}$ Higgs. Since the only configurations of interest are those where the b quarks in the final state have substantial transverse momentum p_T, we ignore the b mass when computing Dirac traces. Except for coupling factors, the matrix elements for CP–even and CP–odd Higgs boson production are then exactly the same. The contribution from $q\bar{q}$ annihilation is found to amount to less than 5% of the contribution from gluon fusion even at the highest Higgs masses accessible at the Tevatron. This contribution can therefore be safely neglected.

In the region of the parameter space of interest to us ($m_A \leq 200$ GeV, $\tan\beta \gg 1$) the decays of A and its nearly mass degenerate CP–even partner are completely dominated by $b\bar{b}$ and $\tau^+\tau^-$ final states [10], which occur with branching ratios of $\simeq 90\%$ and 10%, respectively. Most $b\bar{b}$ Higgs events therefore lead to final states containing four b (anti)quarks. Unfortunately this signal suffers from very large QCD backgrounds; these can be controlled by severe cuts, but then one needs an integrated luminosity of several fb^{-1} for a measurable signal [11].

We therefore focus on Higgs boson decays to $\tau^+\tau^-$ pairs, which lead to far cleaner $b\bar{b}\tau^+\tau^-$ final states.[1] Indeed, the CDF collaboration has searched for an excess in precisely this final state, in the context of searching for third generation leptoquarks [12]. Thus we can derive bounds on associated $b\bar{b}$ Higgs production by simply applying their cuts to our events, and checking whether the resulting cross section falls above the CDF limit.

τ leptons can decay either purely leptonically, with a branching ratio $B(\tau^- \to l^-\bar{\nu}_l\nu_\tau) \simeq 18\%$ for each mode $l = e,\ \mu$, or into low–multiplicity hadronic states and a ν_τ, with a branching ratio $\simeq 64\%$. More than 95% of all hadronic τ decays produce final states with either one or three long lived charged particles, mostly $\pi^\pm$ (the so–called 1–prong and 3–prong decays). The most likely combination for a $\tau^+\tau^-$ pair is therefore a leptonic decay of one of the τ leptons, while the other decays hadronically; the combined branching ratio for this final state is about 45%. The presence of an isolated electron or muon also allows one to trigger on such events with a high efficiency, assuming that this lepton is sufficiently energetic. This motivated the CDF collaboration to focus on this "lepton + $\tau-$jet" signature. They also require the event to contain at least two additional jets, in order to reduce Drell–Yan backgrounds; however, they do not require these jets to be identified as $b-$jets.

The total list of relevant cuts applied in the CDF analysis is

$$\begin{array}{rl} (i) & p_T(l) > 20 \text{ GeV}, \quad |\eta(l)| < 1.0 \\ (ii) & p_T(\tau - \text{jet}) > 15 \text{ GeV}, \quad |\eta(\tau - \text{jet})| < 1.0 \\ (iii) & \Delta\Phi\left(\vec{p}_T(l),\, \vec{\not{p}}_T\right) < 50^\circ \\ (iv) & \geq 2 \text{ jets with } p_T(\text{jet}) > 10 \text{ GeV} \end{array}$$

Here, η denotes the pseudo–rapidity, and $\Delta\Phi$ the transverse angle between the charged lepton and the missing p_T vector. The first cut ensures that the event can be triggered. Cut (ii) is necessary to suppress QCD backgrounds from jets faking τ's; low p_T QCD jets tend to have low multiplicity, and are thus more likely to look like $\tau-$jets. Cut (iii) has been introduced to suppress backgrounds from W+jets events, where the W decays into an electron or muon and one of the jets fakes a τ; in such events the charged lepton and missing p_T (from the single neutrino in the event) tend to be back–to–back in the transverse plane. Finally, cut (iv) reduces backgrounds from the Drell–Yan production of $\tau^+\tau^-$ pairs.[2] After applying these cuts, CDF finds one candidate event, compared to an estimated background from SM sources

[1]We note in passing that this makes our result nearly independent of the numerical value of the running b mass. The reason is that the total production cross section scales like m_b^2, while the branching ratio into $\tau^+\tau^-$ pairs is given by $m_\tau^2/\left(3m_b^2 + m_\tau^2\right) \simeq \frac{1}{3}m_\tau^2 m_b^{-2}\left(1 - \frac{1}{3}m_\tau^2 m_b^{-2}\right)$; the leading m_b dependence therefore cancels in $\sigma \cdot B$.

[2]In addition, we have applied the CDF isolation cut for the lepton, requiring the total transverse energy in a cone with radius $R \equiv \sqrt{(\delta\eta)^2 + (\delta\phi)^2} = 0.4$ around the lepton momentum to be less than 10% of the transverse energy of the lepton. However, this cut has almost no effect in our parton–level Monte Carlo simulation.

of about 2.5 events. This allows them to place an upper limit of 38 fb (at 95% c.l.) on any non–SM contribution to the final state defined by these cuts.

In our simulation we generate $b\bar{b}$ Higgs events and let the Higgs boson decay into $\tau^+\tau^-$ pairs. We model hadronic τ decays as a sum of two–body decays into $\pi\nu_\tau$, $\rho\nu_\tau$ and $a_1\nu_\tau$, with branching ratios given in the literature [13]. We next apply the cuts $(i) - (iv)$. Note that we do not include any out–of–cone losses for the $b-$jets, which would reduce the cross section; nor do we attempt to model energy smearing or initial state radiation, which increase the cross section. For example, some 30% of all $t\bar{t}$ events produce at least one 10 GeV jet from initial state radiation; this fraction could be even higher in our case, since the initial state contains gluons rather than quarks. We therefore consider our estimate of the cross section to be quite conservative. Finally, the event weight is multiplied with a factor describing the τ jet identification efficiency. CDF quotes values of 32% for $p_T(\tau - \text{jet}) \leq 20$ GeV, and 59% for $p_T(\tau - \text{jet}) \geq 40$ GeV; we use linear interpolation for intermediate values of $p_T(\tau - \text{jet})$.

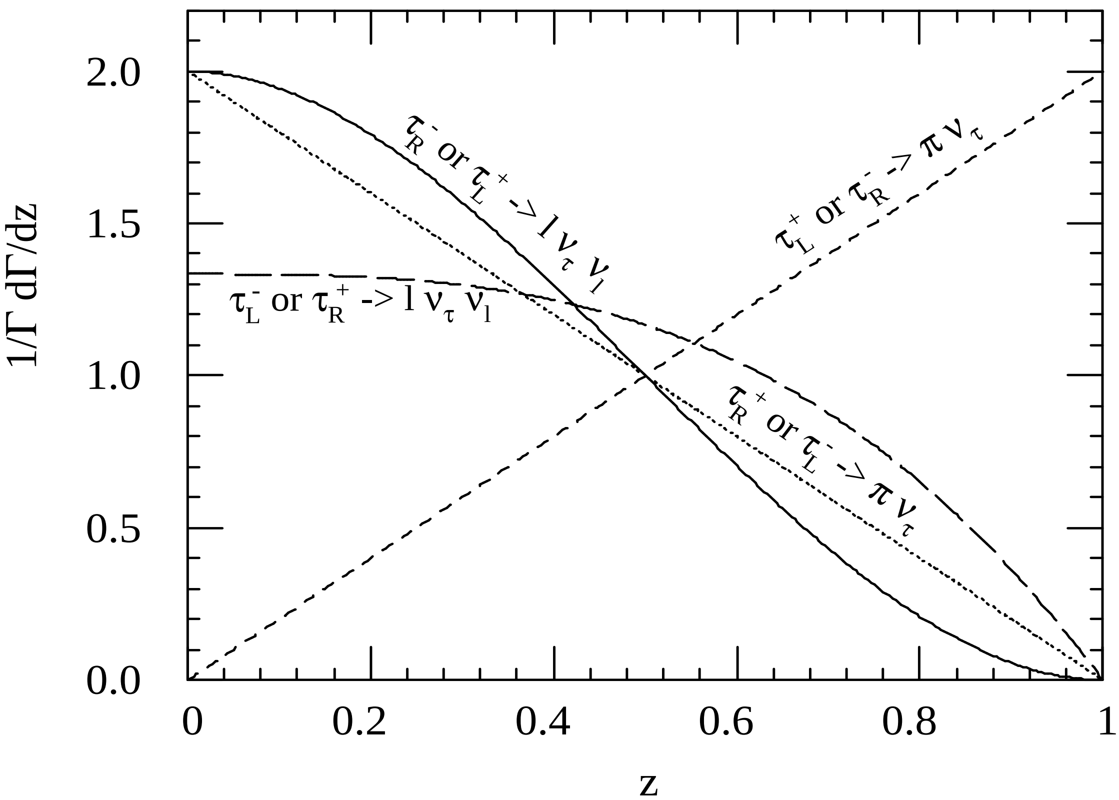

Figure 1: The energy spectrum of the visible τ decay products for τ decays into leptons or into pions for both polarization states of the τ lepton, as a function of $z = E_{l,\pi}/E_\tau$. The τ lepton is assumed to be ultra–relativistic, so that terms $\mathcal{O}(m_\tau/E_\tau)$ can be neglected. In Higgs $\to \tau^+\tau^-$ events, either both τ leptons decay with a "soft" spectrum (solid and dotted curves), or both decay with a "hard" spectrum (short and long dashed curves). In contrast, in the Drell–Yan background events one "soft" decay spectrum is always combined with one "hard" spectrum.

There is one slight subtlety in the treatment of Higgs $\to \tau^+\tau^-$ decays. Even though each τ lepton is unpolarized on the average, the polarizations of the τ^+ and τ^- are completely correlated, since spin–0 bosons can only decay into $\tau_R^+\tau_R^-$ or $\tau_L^+\tau_L^-$ final states. This spin correlation affects the efficiency with which our events pass the cuts $(i)-(iii)$, since the τ decay spectrum depends quite strongly on the τ polarization [14]. This is illustrated in Fig. 1, which shows the energy spectrum of the visible decay products of ultra–relativistic τ leptons decaying either leptonically or into $\pi\nu_\tau$. We can assume without any loss of generality that τ^- decays leptonically while τ^+ decays hadronically. The combination $\tau_R^-\tau_R^+$ then has soft decay spectra for both τ leptons, leading to a very low probability to pass cuts (i), (ii). On the other hand, for Higgs $\to \tau_L^-\tau_L^+$ decays both τ leptons decay with a hard spectrum for the visible decay products, so these events pass cuts (i), (ii) much more easily. Furthermore, the very hard pion spectrum from $\tau_L^+ \to \pi^+\bar{\nu}_\tau$ decays means that the missing p_T mostly comes from the neutrinos produced in the leptonic decays of τ_L^-; most of these events therefore also pass cut (iii). The situation is similar for $\tau^+ \to \rho^+\bar{\nu}_\tau$ decays, although the ρ energy spectrum is less sensitive to the τ polarization than the π spectrum depicted in Fig. 1. Finally, the a_1 spectrum from $\tau^+ \to a_1^+\bar{\nu}_\tau$ decays is almost independent of the τ polarization [14]; cut (i) nevertheless favors the combination $\tau_L^-\tau_L^+$.

In contrast, Drell–Yan backgrounds only produce $\tau_L^-\tau_R^+$ and $\tau_R^-\tau_L^+$ final states. This means that one τ lepton will decay with a hard spectrum, and the other one with a soft spectrum of the visible decay products. In case of $\tau^+ \to \pi^+$ or ρ^+ decays, the efficiency of passing either cut (i) or cut (ii) is therefore always quite low. As a result, the total efficiency of Drell–Yan background events in the $l\pi$ channel is nearly a factor of 2 lower than what it is for signal events. The difference is less pronounced for $l\rho$ or la_1 final states, but the overall efficiency of background events remains lower also in these channels. This enhancement of the signal to background ratio is no less welcome for being accidental; note that CDF searched [12] for $\tau^+\tau^-$ pairs without any spin correlation.

As emphasized earlier, there is no evidence for any non–SM contribution to the final state defined by cuts $(i)-(iv)$, which allowed the CDF collaboration to place an upper bound of 38 fb on any such contribution. Within the framework of the MSSM this translates into the m_A dependent upper bound on $\tan\beta$ shown by the solid line in Fig. 2. Recall that either h or H contributes equally to the $b\bar{b}\tau^+\tau^-$ final state for $\tan\beta \gg 1$. This is not necessarily true for general two Higgs doublet models (2HDM); in these models the resulting bound on $\tan\beta$ will therefore be weaker by up to a factor of $\sqrt{2}$. We have also indicated the present LEP bound $m_A > 75$ GeV, as well as the bound that can be deduced from an analysis [6] of top decays (dashed curve). Both these bounds are valid only in the MSSM.[3]

The new bound shown in Fig. 2 has far reaching ramifications. For lack of space we merely mention two examples. First, simple $SO(10)$ models require $\tan\beta \simeq 50$ in order to unify all three third generation Yukawa couplings [15]. For $\tan\beta = 50$, our bound excludes the region $m_A \leq 190$ GeV. This excludes the entire branch of "low m_A" solutions found by Blazek and Raby in their recent analysis of a complete $SO(10)$ model [16].

Furthermore, our new constraint will reduce the maximal expected event rate in laboratory experiments searching for relic neutralinos left over from the Big Bang; these are one of the best

[3]Of course, LEP searches, top decays, and radiative flavor changing decays such as $b \to s\gamma$ do constrain general 2HDM; however, these constraints do not lead to bounds on m_A that are valid for all choices of the other parameters of these models.

candidates for the mysterious "Dark Matter" in the Universe [17]. In particular, it excludes the scenarios found by Bergström et al. [18]. which might have led to event rates within the reach of current experiments. These scenarios require low m_A and large tanβ to maximize the Higgs exchange contribution to the neutralino–nucleon scattering cross section.

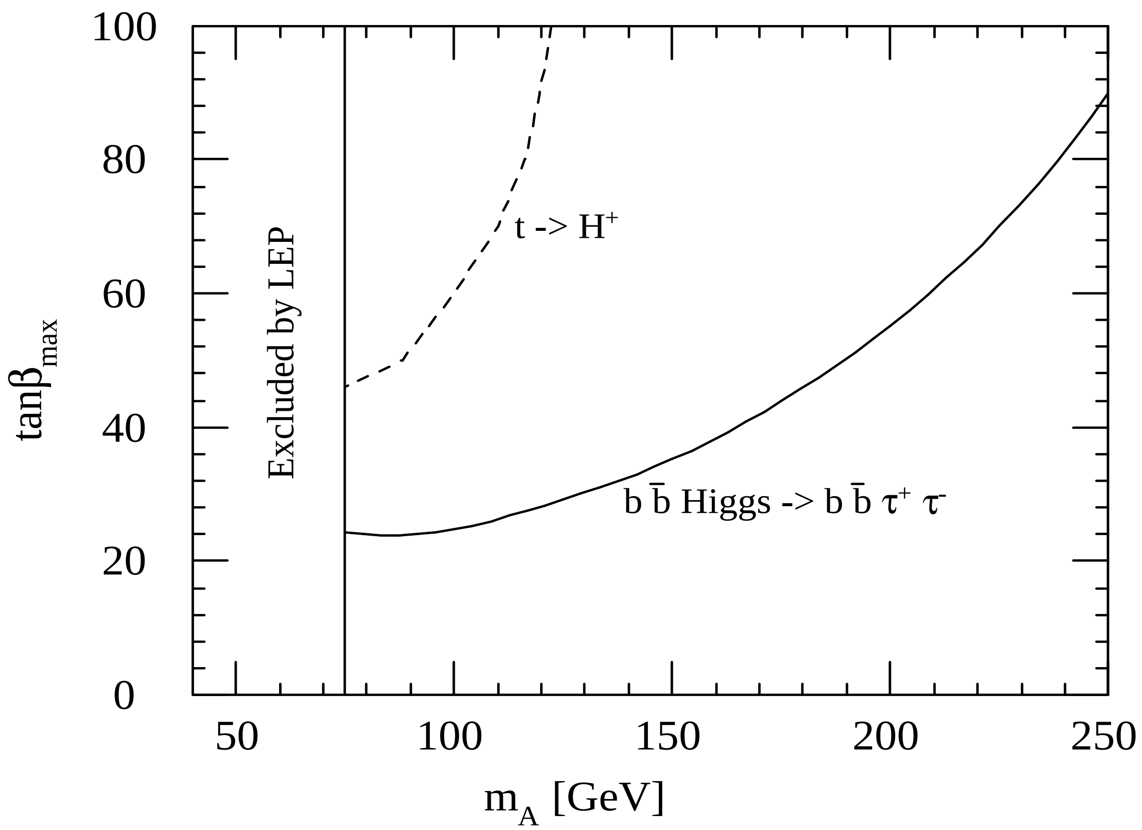

Figure 2: Constraints on the MSSM Higgs sector in the (m_A, tanβ) plane. The region $m_A < 75$ GeV is excluded by Higgs searches at the LEP collider. The region above the dashed curve is excluded by an analysis [6] of top quark decays. The region above the solid line is excluded by our analysis using CDF limits on $b\bar{b}\tau^+\tau^-$ final states. For tan$\beta \gg 1$, the CP–odd scalar A is nearly degenerate with either the light CP–even scalar h or the heavy CP–even state H; the cross–over occurs between 100 and 125 GeV, depending on details of the sparticle spectrum.

In summary, we have exploited results on $\tau^+\tau^- + 2$ jets final states published by the CDF collaboration to place stringent new limits on the Higgs sector of the Minimal Supersymmetric Standard Model in the theoretically interesting region of large tanβ. These bounds are valid in a general MSSM, independent of the details of the sparticle spectrum. If the mass of the CP–odd Higgs boson is just beyond the present LEP bound, the region tan$\beta \geq 25$ is excluded. Conversely, if tan$\beta = 50$, our bound extends to $m_A \simeq 190$ GeV. More detailed studies of the $b\bar{b}\tau^+\tau^-$ final state as a means to discover MSSM Higgs bosons at future runs of the Tevatron collider or at the LHC will be presented elsewhere.

Acknowledgements

We have gained from discussions with D.P. Roy and N.K. Mondal. M.D. thanks the Theoretical Physics Department at TIFR for their kind hospitality. M.G. was supported in part by a 1997 foreign postdoctoral fellowship through the Korean Science and Engineering Foundation, and in part by a graduate school fellowship from Yonsei University.

References

[1] For a review, see e.g. H.P. Nilles, Phys. Rep. **110**, 1 (1984).

[2] P. Fayet, Phys. Lett. **64B**, 159 (1976).

[3] J.F. Gunion and H.E. Haber, Nucl. Phys. **B272**, 1 (1986).

[4] Y. Okada, M. Yamaguchi and T. Yanagida, Prog. Theor. Phys. **85**, 1 (1991), and Phys. Lett. **B262**, 54 (1991); H.E. Haber and R. Hempfling, Phys. Rev. Lett. **66**, 1815 (1991); J. Ellis, G. Ridolfi and F. Zwirner, Phys. Lett. **B257**, 83 (1991).

[5] P. Janot, talk at *SUSYUNI: International Conference on Supersymmetry and Unification*, Mumbai, India, December 1997.

[6] M. Guchait and D.P. Roy, Phys. Rev. **D55**, 7263 (1997). See also CDF collab., F. Abe et al., Phys. Rev. Lett. **79**, 357 (1997).

[7] CDF collab., F. Abe et al., FERMILAB-PUB-97/286-E; D0. collab., S. Abachi et al., Phys. Rev. Lett. **79**, 1203 (1997).

[8] A. Stange, W. Marciano and S. Willenbrock, Phys. Rev. **D49**, 1354 (1994); S. Mrenna and G.L. Kane, hep–ph/9406337.

[9] J. N. Ng and P. Zakarauskas, Phys. Rev. **D29**, 876 (1984).

[10] J.F. Gunion, H.E. Haber, G.L. Kane and S. Dawson. *The Higgs Hunter's Guide*, Addison–Wesley, 1992.

[11] J. Dai, J.F. Gunion and R. Vega, Phys. Lett. **B387**, 801 (1996).

[12] CDF collab., F. Abe et al., Phys. Rev. Lett. **78**, 2906 (1997).

[13] Particle Data Group, R.M. Barnett et al., Phys. Rev. **D54**, 1 (1996).

[14] B.K. Bullock, K. Hagiwara and A.D. Martin, Nucl. Phys. **B395**, 499 (1993).

[15] B. Ananthanarayan, G. Lazarides and Q. Shafi, Phys. Rev. **D44**, 1613 (1991).

[16] T. Blazek and S. Raby, hep–ph/9712257.

[17] For a review, see G. Jungman, M. Kamionkowski, and K. Griest, Phys. Rep. **267**, 195 (1996).

[18] L. Bergström and P. Gondolo, Astropart. Phys. **5**, 263 (1995).