Pre-LHC SUSY Searches: an Overview[*]

A. Masiero

SISSA, Via Beirut 2-4, 34013 Trieste, Italy and
INFN, sez. di Trieste, Padriciano 99, 34012 Trieste,Italy

and

L. Silvestrini

Technische Universität München, Physik Department,
D-85748 Garching, Germany

Abstract

We discuss the prospects for searches of low-energy supersymmetry in the time interval separating us from the advent of LHC. In this period of time "indirect" searches may play a very relevant role. We refer to manifestations of supersymmetry in flavour changing neutral current and CP violating phenomena and to signals of the lightest supersymmetric particle in searches of dark matter. In the first part of the talk we critically review the status of the minimal supersymmetric model to discuss the chances that direct and indirect supersymmetric searches may have before the LHC start. In the second part we point out what we consider to be the most promising grounds where departures from the standard model prediction may signal the presence of new physics, possibly of supersymmetric nature. We argue that the often invoked complementarity of direct and indirect searches of low-energy supersymmetry is becoming even more true in the pre-LHC era.

1 Introduction

It is not rare to hear the following gloomy forecast: if no supersymmetric signal is seen at LEP, then we have nothing else to do but wait for LHC. We do not agree with this statement. Apart from the fact that even for direct searches one should take into account the relevant potentialities of Tevatron in the LEP-LHC time interval, one should not neglect that indirect searches for new physics signals are going to be flourishing before 2005. We refer to processes exploring flavour physics (with or without CP violation) where new particles can play an active role being exchanged in the loop contributions and to several new astroparticle observations which may constitute privileged places to obtain information on physics beyond the Standard Model (SM).

We wish to present here a brief overview (which is necessarily biased by our theoretical prejudices) of what we consider most promising in this effort of looking for indirect signals of low-energy Supersymmetry (SUSY) before the LHC advent. First we will review the status and prospects for direct SUSY searches, then we will discuss the role that SUSY may play in Flavour Changing Neutral Current (FCNC) and CP violating phenomena. Finally we will briefly comment on searches for the lightest SUSY particle in experiments looking for Dark Matter (DM).

2 Status of the MSSM

It is known that, even asking for the minimal content of superfields which are necessary to supersymmetrize the SM and imposing R parity, one is still left with more than 100 free parameters most of which are in the flavour sector. It is also true that very large portions on this enormous SUSY parameter space are already ruled out by the present phenomenology (in particular FCNC and CP constraints). If one wants to reduce the number of free parameters one has to make assumptions on what lies well beyond low-energy SUSY, in particular on the quite unknown issue of the origin of SUSY breaking. The two most popular drastic reductions of free SUSY parameters are provided by minimal supergravity (SUGRA) [1] (with the further assumption of unification of gauge couplings and gaugino masses at some grand unification scale) and by the models of Gauge-Mediated SUSY Breaking (GMSB) [2]-[4]. In minimal SUGRA and the minimal version of GMSB we have only 3 or 4 parameters

in addition to those of the SM and so we can become much more predictive.

In the context of the minimal supergravity model (with electroweak radiative breaking), we ask the following relevant questions for direct SUSY searches: i) given the present experimental lower bounds on the masses of SUSY particles, how much room have we got in the SUSY parameter space to explore, or, in other words, when should we give up with SUSY if searches are fruitless? ii) is there any experimental signature of low-energy SUSY which is independent from the choice of the SUSY parameters in particular of the soft breaking sector? iii) are the electroweak precision tests telling us something relevant on low-energy SUSY?

i) SUSY must be a low-energy symmetry if it has to deal with the issue of the gauge hierarchy problem. This fact is usually translated into the statement that SUSY particle masses should not be significantly larger than $O(1\,\text{TeV})$ given that SUSY breaking should not exceed this energy scale to realize a suitable "protection" of the mass of the scalar Higgs responsible for the electroweak breaking. Actually one may try to be more quantitative [5]. First one relates the Z mass to the value of the 4 parameters of the minimal SUGRA run from the large scale, at which the soft breaking terms originate, down to the electroweak scale. Then one establishes a degree of naturalness corresponding to the amount of fine tuning of the initial SUSY parameters which is needed to reproduce the correct Z mass for increasing values of the low-energy SUSY masses. For instance it is clear that to have all SUSY particles with a mass of $O(1\,\text{TeV})$ would require a severe fine tuning of the boundary conditions.

As for all naturalness criteria, also in this case there is a large amount of subjectivity, but one message emerges quite clearly: already now, in particular with the lower bound on chargino masses exceeding 90 GeV, we are entering an area of parameter space where a certain degree of fine tuning is needed. Hence we are already at the stage where we may expect "naturally" to find SUSY particles. Moreover such naturalness analyses confirm that LHC represents kind of "definitive" machine for SUSY direct searches: if no SUSY particle is discovered at LHC, the degree of fine tuning becomes so severe that it is hard to still defend the idea of low-energy SUSY. Finally, an important comment on the degree at which different SUSY masses are constrained by such naturalness criteria: due to the large difference in the Yukawa couplings of the

2

third (heaviest) generation with respect to the first two generations, it turns out that only the sbottoms and stops are required not to be very heavy, whilst squarks of the first two generations can be quite heavy (say tens of TeV) without severely affecting the correct electroweak breaking. This observations may play a relevant role in tackling the FCNC problem in SUSY (see below).

ii) If one allows the SUSY parameters to take larger and larger values, all the SUSY particles become heavier and heavier with only one remarkable exception. In the Higgs mass spectrum of SUSY models the lightest scalar always remains light. The mass of the light CP-even neutral Higgs in the MSSM is calculable at tree level in terms of two SUSY parameters of the Higgs potential. At this level it is smaller than the mass of the Z. When radiative corrections are included, the mass of the light Higgs becomes a function also of the other SUSY parameters and its upper bound increases significantly [6]. However, even varying the MSSM parameters as much as one wishes, it is not possible to exceed $130 - 135$ GeV for its mass. Indeed, taking $m_t = 175$ GeV and for a stop lighter than 1 TeV one obtains that the upper bound on the lightest Higgs is 125 GeV allowing for "maximal" mixing in the top squark sector (the bound decreases for smaller stop mixing).

It is not easy to significantly evade the above upper bound on the mass of the lightest Higgs even if one gives up the minimality of the SUSY model. For instance, if one adds a singlet to the two Higgs doublets (i.e., one goes to the so called Next-to-Minimal SUSY Standard Model, NMSSM), then a new parameter shows up in the scalar potential: the coupling of the singlet with the two doublets. If one imposes that all couplings remain perturbative up to the Planck scale, then the consequent upper bound on this new coupling implies that the lightest Higgs should not be heavier than 150 GeV or so [7].

Obviously having a possibly "exotic" Higgs below 150 GeV does not necessarily mean that it can be seen at LHC. While the lightest Higgs of the MSSM seems to be detectable at LHC, there may still be some significant loopholes for searches of the light Higgs in the NMSSM context.

iii) It is known that the MSSM is a decoupling theory. In the limit where

3

we send the SUSY parameters to infinity all SUSY masses become infinite, with only the lightest Higgs remaining light and coinciding with the usual SM Higgs. In this limit we would recover the SM. It turns out that as far as electroweak precision tests are concerned, the decoupling of the MSSM is quite fast: already for SUSY masses above $200 - 300$ GeV the effects due to the exchange of SUSY particles in radiative corrections to the electroweak observables become negligible. Notice that this is not true if, instead of electroweak precision tests, we consider FCNC and CP tests. In this latter case, the decoupling may be much slower with squarks and gluinos of 1 TeV still providing sizeable contributions in loop diagrams to some rare processes.

Obviously the SM fit of electroweak precision data is now so good that there is no point in trying to improve it by the addition of the several degrees of freedom represented by the SUSY particles. The situation was different a couple of years ago when the discrepancy between the SM prediction and the data in the decay of the Z into a b quark pair resulted in a SM fit which could be significantly improved. Now the goal of the game has changed: one looks for regions of the SUSY parameter space where (some) SUSY masses are sufficiently small so that virtual SUSY contributions to electroweak observables are sizeable [8]. Some of these regions may cause unbearably high departures from the SM predictions and hence they can be ruled out. In this way it is possible to exclude some (limited) portions of the MSSM parameter space which would be otherwise allowed by the limits on SUSY parameters coming from direct searches of SUSY particles.

Finally, we make a comment related to the prediction of one low-energy parameter (the electroweak angle or, as it is the case nowadays, the value of the strong coupling at the Z mass scale) when one asks for the unification of the gauge coupling constants in the MSSM. The value predicted for $\alpha_S(m_Z)$ in the MSSM is a couple of standard deviations higher than the experimental value. We do not consider this as a problem for the MSSM. Indeed, high energy thresholds generated from the masses of superheavy GUT particles may conceivably produce corrections able to account for such discrepancy. Taking into account the uncertainties in the dynamics at the GUT scale, we consider the argument of unification of couplings as a support to the existence of

low-energy SUSY.

Before starting our discussion of indirect searches of SUSY, let us emphasise that direct production and detection of SUSY particles remain the only way to definitely prove the existence of low-energy SUSY. However it is true that if LEP II is not going to find a SUSY signal and unless some surprise possibly comes from Tevatron, we will have to wait almost ten years to obtain an answer from such direct searches. In view of this fact and of what we said in this section we think that indirect searches of SUSY in the pre-LHC era deserve a very special attention.

3 FCNC and SUSY

The generation of fermion masses and mixings ("flavour problem") gives rise to a first and important distinction among theories of new physics beyond the electroweak standard model.

One may conceive a kind of new physics which is completely "flavour blind", i.e. new interactions which have nothing to do with the flavour structure. To provide an example of such a situation, consider a scheme where flavour arises at a very large scale (for instance the Planck mass) while new physics is represented by a supersymmetric extension of the SM with supersymmetry broken at a much lower scale and with the SUSY breaking transmitted to the observable sector by flavour-blind gauge interactions [2]-[4]. In this case one may think that new physics does not cause any major change to the original flavour structure of the SM, namely that the pattern of fermion masses and mixings is compatible with the numerous and demanding tests of flavour changing neutral currents.

Alternatively, one can conceive a new physics which is entangled with the flavour problem. As an example consider a technicolour scheme where fermion masses and mixings arise through the exchange of new gauge bosons which mix together ordinary and technifermions. Here we expect (correctly enough) new physics to have potential problems in accommodating the usual fermion spectrum with the adequate suppression of FCNC. As another example of new physics which is not flavour blind, take a more conventional SUSY model which is derived from a spontaneously broken N=1 supergravity and where the SUSY breaking information is conveyed to the ordinary sector of the theory through gravitational interactions. In this case we may

expect that the scale at which flavour arises and the scale of SUSY breaking are not so different and possibly the mechanism itself of SUSY breaking and transmission is flavour-dependent. Under these circumstances we may expect a potential flavour problem to arise, namely that SUSY contributions to FCNC processes are too large.

The potentiality of probing SUSY in FCNC phenomena was readily realized when the era of SUSY phenomenology started in the early 80's [9]. In particular, the major implication that the scalar partners of quarks of the same electric charge but belonging to different generations had to share a remarkably high mass degeneracy was emphasised.

Throughout the large amount of work in this last decade it became clearer and clearer that generically talking of the implications of low-energy SUSY on FCNC may be rather misleading. In minimal SUGRA FCNC contributions can be computed in terms of a very limited set of unknown new SUSY parameters. Remarkably enough, this minimal model succeeds to pass all the set of FCNC tests unscathed. To be sure, it is possible to severely constrain the SUSY parameter space, for instance using $b \to s\gamma$, in a way which is complementary to what is achieved by direct SUSY searches at colliders.

However, the MSSM is by no means equivalent to low-energy SUSY. A first sharp distinction concerns the mechanism of SUSY breaking and transmission to the observable sector which is chosen. As we mentioned above, in models with gauge-mediated SUSY breaking (GMSB models [2]-[4]) it may be possible to avoid the FCNC threat "ab initio" (notice that this is not an automatic feature of this class of models, but it depends on the specific choice of the sector which transmits the SUSY breaking information, the so-called messenger sector). The other more "canonical" class of SUSY theories that was mentioned above has gravitational messengers and a very large scale at which SUSY breaking occurs. In this talk we will focus only on this class of gravity-mediated SUSY breaking models. Even sticking to this more limited choice we have a variety of options with very different implications for the flavour problem.

First, there exists an interesting large class of SUSY realizations where the customary R-parity (which is invoked to suppress proton decay) is replaced by other discrete symmetries which allow either baryon or lepton violating terms in the superpotential. But, even sticking to the more orthodox view of imposing R-parity, we are still left with a large variety of extensions of the MSSM at low energy. The point is that low-energy SUSY "feels" the new

physics at the superlarge scale at which supergravity (i.e., local supersymmetry) broke down. In this last couple of years we have witnessed an increasing interest in supergravity realizations without the so-called flavour universality of the terms which break SUSY explicitly. Another class of low-energy SUSY realizations which differ from the MSSM in the FCNC sector is obtained from SUSY-GUT's. The interactions involving superheavy particles in the energy range between the GUT and the Planck scale bear important implications for the amount and kind of FCNC that we expect at low energy.

Given a specific SUSY model it is in principle possible to make a full computation of all the FCNC phenomena in that context. However, given the variety of options for low-energy SUSY (even confining ourselves here to models with R matter parity), it is important to have a way to extract from the whole host of FCNC processes a set of upper limits on quantities which can be readily computed in any chosen SUSY frame.

The best model-independent parameterisation of FCNC effects is the so-called mass insertion approximation [10]. It concerns the most peculiar source of FCNC SUSY contributions that do not arise from the mere supersymmetrization of the FCNC in the SM. They originate from the FC couplings of gluinos and neutralinos to fermions and sfermions [11]. One chooses a basis for the fermion and sfermion states where all the couplings of these particles to neutral gauginos are flavour diagonal, while the FC is exhibited by the non-diagonality of the sfermion propagators. Denoting by Δ the off-diagonal terms in the sfermion mass matrices (i.e. the mass terms relating sfermion of the same electric charge, but different flavour), the sfermion propagators can be expanded as a series in terms of $\delta = \Delta/\tilde{m}^2$, where $\tilde{m}$ is the average sfermion mass. As long as Δ is significantly smaller than $\tilde{m}^2$, we can just take the first term of this expansion and, then, the experimental information concerning FCNC and CP violating phenomena translates into upper bounds on these δ's [12]-[14].

Obviously the above mass insertion method presents the major advantage that one does not need the full diagonalisation of the sfermion mass matrices to perform a test of the SUSY model under consideration in the FCNC sector. It is enough to compute ratios of the off-diagonal over the diagonal entries of the sfermion mass matrices and compare the results with the general bounds on the δ's that we provide here from all available experimental information.

There exist four different Δ mass insertions connecting flavours i and j along a sfermion propagator: $(\Delta_{ij})_{LL}$, $(\Delta_{ij})_{RR}$, $(\Delta_{ij})_{LR}$ and $(\Delta_{ij})_{RL}$. The

x	$\sqrt{\left\|\mathrm{Re}\left(\delta_{12}^d\right)_{LL}^2\right\|}$	$\sqrt{\left\|\mathrm{Re}\left(\delta_{12}^d\right)_{LR}^2\right\|}$	$\sqrt{\left\|\mathrm{Re}\left(\delta_{12}^d\right)_{LL}\left(\delta_{12}^d\right)_{RR}\right\|}$
0.3	1.9×10^{-2}	7.9×10^{-3}	2.5×10^{-3}
1.0	4.0×10^{-2}	4.4×10^{-3}	2.8×10^{-3}
4.0	9.3×10^{-2}	5.3×10^{-3}	4.0×10^{-3}

x	$\sqrt{\left\|\mathrm{Re}\left(\delta_{13}^d\right)_{LL}^2\right\|}$	$\sqrt{\left\|\mathrm{Re}\left(\delta_{13}^d\right)_{LR}^2\right\|}$	$\sqrt{\left\|\mathrm{Re}\left(\delta_{13}^d\right)_{LL}\left(\delta_{13}^d\right)_{RR}\right\|}$
0.3	4.6×10^{-2}	5.6×10^{-2}	1.6×10^{-2}
1.0	9.8×10^{-2}	3.3×10^{-2}	1.8×10^{-2}
4.0	2.3×10^{-1}	3.6×10^{-2}	2.5×10^{-2}

x	$\sqrt{\left\|\mathrm{Re}\left(\delta_{12}^u\right)_{LL}^2\right\|}$	$\sqrt{\left\|\mathrm{Re}\left(\delta_{12}^u\right)_{LR}^2\right\|}$	$\sqrt{\left\|\mathrm{Re}\left(\delta_{12}^u\right)_{LL}\left(\delta_{12}^u\right)_{RR}\right\|}$
0.3	4.7×10^{-2}	6.3×10^{-2}	1.6×10^{-2}
1.0	1.0×10^{-1}	3.1×10^{-2}	1.7×10^{-2}
4.0	2.4×10^{-1}	3.5×10^{-2}	2.5×10^{-2}

Table 1: Limits on $\mathrm{Re}\left(\delta_{ij}\right)_{AB}\left(\delta_{ij}\right)_{CD}$, with $A, B, C, D = (L, R)$, for an average squark mass $m_{\tilde{q}} = 500\mathrm{GeV}$ and for different values of $x = m_{\tilde{g}}^2/m_{\tilde{q}}^2$. For different values of $m_{\tilde{q}}$, the limits can be obtained multiplying the ones in the table by $m_{\tilde{q}}(\mathrm{GeV})/500$.

indices L and R refer to the helicity of the fermion partners. Instead of the dimensional quantities Δ it is more useful to provide bounds making use of dimensionless quantities, δ, that are obtained dividing the mass insertions by an average sfermion mass.

Let us first consider CP-conserving $\Delta F = 2$ processes. The amplitudes for gluino-mediated contributions to $\Delta F = 2$ transitions in the mass-insertion approximation have been computed in refs. [13, 14]. Imposing that the contribution to $K - \bar{K}$, $D - \bar{D}$ and $B_d - \bar{B}_d$ mixing proportional to each single δ parameter does not exceed the experimental value, we obtain the constraints on the δ's reported in table 1, barring accidental cancellations [14] (for a QCD-improved computation of the constraints coming from $K - \bar{K}$ mixing, see ref. [15]).

We then consider the process $b \rightarrow s\gamma$. This decay requires a helicity

flip. In the presence of a $\left(\delta_{23}^d\right)_{LR}$ mass insertion we can realize this flip in the gluino running in the loop. On the contrary, the $\left(\delta_{23}^d\right)_{LL}$ insertion requires the helicity flip to occur in the external b-quark line. Hence we expect a stronger bound on the $\left(\delta_{23}^d\right)_{LR}$ quantity. Indeed, this is what happens: $\left(\delta_{23}^d\right)_{LL}$ is essentially not bounded, while $\left(\delta_{23}^d\right)_{LR}$ is limited to be $< 10^{-3} - 10^{-2}$ according to the average squark and gluino masses [14].

Given the upper bound on $\left(\delta_{23}^d\right)_{LR}$ from $b \to s\gamma$, it turns out that the quantity x_s of the $B_s - \bar{B}_s$ mixing receives contributions from this kind of mass insertions which are very tiny. The only chance to obtain large values of x_s is if $\left(\delta_{23}^d\right)_{LL}$ is large, say of $O(1)$. In that case x_s can easily jump up to values of $O(10^2)$ or even larger.

Then, imposing the bounds in table 1, we can obtain the largest possible value for $\mathrm{BR}(b \to d\gamma)$ through gluino exchange. As expected, the $\left(\delta_{13}^d\right)_{LL}$ insertion leads to very small values of this BR of $O(10^{-7})$ or so, whilst the $\left(\delta_{13}^d\right)_{LR}$ insertion allows for $\mathrm{BR}(b \to d\gamma)$ ranging from few times 10^{-4} up to few times 10^{-3} for decreasing values of $x = m_{\tilde{g}}^2/m_{\tilde{q}}^2$. In the SM we expect $\mathrm{BR}(b \to d\gamma)$ to be typically $10 - 20$ times smaller than $\mathrm{BR}(b \to s\gamma)$, i.e. $\mathrm{BR}(b \to d\gamma) = (1.7 \pm 0.85) \times 10^{-5}$. Hence a large enhancement in the SUSY case is conceivable if $\left(\delta_{13}^d\right)_{LR}$ is in the 10^{-2} range. Notice that in the MSSM we expect $\left(\delta_{13}^d\right)_{LR} < m_b^2/m_{\tilde{q}}^2 \times V_{td} < 10^{-6}$, hence with no hope at all of a sizeable contribution to $b \to d\gamma$.

An analysis similar to the one of $b \to s\gamma$ decays can be performed in the leptonic sector where the masses $m_{\tilde{q}}$ and $m_{\tilde{g}}$ are replaced by the average slepton mass $m_{\tilde{l}}$ and the photino mass $m_{\tilde{\gamma}}$ respectively. The most stringent bound concerns the transition $\mu \to e\gamma$ with $\left(\delta_{12}^l\right)_{LR} < 10^{-6}$ for slepton and photino masses of $O(100\ \mathrm{GeV})$ [14].

4 CP and SUSY

The situation concerning CP violation in the MSSM case with $\Phi_A = \Phi_B = 0$ and exact universality in the soft-breaking sector can be summarised in the following way: the MSSM does not lead to any significant deviation from the SM expectation for CP-violating phenomena as d_N^e, ε, ε' and CP violation

9

in B physics; the only exception to this statement concerns a small portion of the MSSM parameter space where a very light $\tilde{t}$ ($m_{\tilde{t}} < 100$ GeV) and χ^+ ($m_\chi \sim 90$ GeV) are present. In this latter particular situation sizeable SUSY contributions to ε_K are possible and, consequently, major restrictions in the $\rho - \eta$ plane can be inferred (see, for instance, ref. [16]). Obviously, CP violation in B physics becomes a crucial test for this MSSM case with very light $\tilde{t}$ and χ^+. Interestingly enough, such low values of SUSY masses are at the border of the detectability region at LEP II.

We now turn to CP violation in the model-independent approach that we are proposing here. For a detailed discussion we refer the reader to our general study [14]. Here we just summarise the situation in the following three points:

i) ϵ provides bounds on the imaginary parts of the quantities whose real part was limited by the K mass difference which are roughly one order of magnitude more severe than the corresponding ones derived from Δm_K.

ii) The nature of the SUSY contribution to CP violation is generally super-weak, since the constraints from ε are always stronger (in the left-left sector) or at least equal (in the left-right sector) to the ones coming from ε'/ε.

iii) the experimental bound on the electric dipole moment of the neutron imposes very stringent limits on $\mathrm{Im}\left(\delta_{11}^d\right)_{LR}$ (of $O(10^{-6})$ for an average squark and gluino mass of 500 GeV.) In conclusion, although technically it is conceivable that some SUSY extension may provide a sizable contribution to ε'/ε, it is rather difficult to imagine how to reconcile a relatively large value of $\mathrm{Im}\left(\delta_{12}^d\right)_{LR}$ with the very strong constraint on the flavour-conserving $\mathrm{Im}\left(\delta_{11}^d\right)_{LR}$ from d_N^e.

We now move to the next frontier for testing the unitarity triangle in general and in particular CP violation in the SM and its SUSY extensions: B physics. We have seen above that the transitions between 1st and 2nd generation in the down sector put severe constraints on $\mathrm{Re}\,\delta_{12}^d$ and $\mathrm{Im}\,\delta_{12}^d$ quantities. To be sure, the bounds derived from ε and ε' are stronger than the corresponding bounds from ΔM_K. If the same pattern repeats itself in

10

the transition between 3rd and 1st or 3rd and 2nd generation in the down sector we may expect that the constraints inferred from $B_d - \bar{B}_d$ oscillations or $b \to s\gamma$ do not prevent conspicuous new contributions also in CP violating processes in B physics. We are going to see below that this is indeed the case ad we will argue that measurements of CP asymmetries in several B-decay channels may allow to disentangle SM and SUSY contributions to the CP decay phase.

New physics can modify the SM predictions on CP asymmetries in B decays by changing the phase of the B_d–$\bar{B}_d$ mixing and the phase and absolute value of the decay amplitude. The general SUSY extension of the SM that we discuss here affects both these quantities.

The crucial question is then: where and how can one possibly distinguish SUSY contributions to CP violation in B decays [17]?

In terms of the decay amplitude A, the CP asymmetry reads

$$\mathcal{A}(t) = \frac{(1 - |\lambda|^2)\cos(\Delta M_d t) - 2\mathrm{Im}\lambda \sin(\Delta M_d t)}{1 + |\lambda|^2} \tag{1}$$

with $\lambda = e^{-2i\phi^M} \bar{A}/A$. In order to be able to discuss the results model-independently, we have labeled as ϕ^M the generic mixing phase. The ideal case occurs when one decay amplitude only appears in (or dominates) a decay process: the CP violating asymmetry is then determined by the total phase $\phi^T = \phi^M + \phi^D$, where ϕ^D is the weak phase of the decay. This ideal situation is spoiled by the presence of several interfering amplitudes.

We summarise the results in table 2 which is taken from the recent analysis of ref. [18]. We refer the interested reader to our work [18] for all the details of how our computation in the SM and in SUSY is carried out. Φ_{SM}^D denotes the decay phase in the SM; for each channel, when two amplitudes with different weak phases are present, we indicate the SM phase of the Penguin (P) and Tree-level (T) decay amplitudes. For $B \to K_S \pi^0$ the penguin contributions (with a vanishing phase) dominate over the tree-level amplitude because the latter is Cabibbo suppressed. For the channel $b \to s\bar{s}d$ only penguin operators or penguin contractions of current-current operators contribute. The phase γ is present in the penguin contractions of the $(\bar{b}u)(\bar{u}d)$ operator, denoted as u-P γ in table 2. $\bar{b}d \to \bar{q}q$ indicates processes occurring via annihilation diagrams which can be measured from the last two channels of table 2. In the case $B \to K^+ K^-$ both current-current and penguin operators contribute. In $B \to D^0 \bar{D}^0$ the contributions from the $(\bar{b}u)(\bar{u}d)$ and

the $(\bar{b}c)(\bar{c}d)$ current-current operators (proportional to the phase γ) tend to cancel out.

SUSY contributes to the decay amplitudes with phases induced by δ_{13} and δ_{23} which we denote as ϕ_{13} and ϕ_{23}. The ratios of A_{SUSY}/A_{SM} for SUSY masses of 250 and 500 GeV are reported in the r_{250} and r_{500} columns of table 2.

We now draw some conclusions from the results of table 2. In the SM, the first six decays measure directly the mixing phase β, up to corrections which, in most of the cases, are expected to be small. These corrections, due to the presence of two amplitudes contributing with different phases, produce uncertainties of $\sim 10\%$ in $B \to K_S \pi^0$, and of $\sim 30\%$ in $B \to D^+ D^-$ and $B \to J/\psi \pi^0$. In spite of the uncertainties, however, there are cases where the SUSY contribution gives rise to significant changes. For example, for SUSY masses of O(250) GeV, SUSY corrections can shift the measured value of the sine of the phase in $B \to \phi K_S$ and in $B \to K_S \pi^0$ decays by an amount of about 70%. For these decays SUSY effects are sizeable even for masses of 500 GeV. In $B \to J/\psi K_S$ and $B \to \phi \pi^0$ decays, SUSY effects are only about 10% but SM uncertainties are negligible. In $B \to K^0 \bar{K}^0$ the larger effect, $\sim 20\%$, is partially covered by the indetermination of about 10% already existing in the SM. Moreover the rate for this channel is expected to be rather small. In $B \to D^+ D^-$ and $B \to K^+ K^-$, SUSY effects are completely obscured by the errors in the estimates of the SM amplitudes. In $B^0 \to D^0_{CP} \pi^0$ the asymmetry is sensitive to the mixing angle ϕ_M only because the decay amplitude is unaffected by SUSY. This result can be used in connection with $B^0 \to K_s \pi^0$, since a difference in the measure of the phase is a manifestation of SUSY effects.

Turning to $B \to \pi\pi$ decays, both the uncertainties in the SM and the SUSY contributions are very large. Here we witness the presence of three independent amplitudes with different phases and of comparable size. The observation of SUSY effects in the $\pi^0 \pi^0$ case is hopeless. The possibility of separating SM and SUSY contributions by using the isospin analysis remains an open possibility which deserves further investigation. For a thorough discussion of the SM uncertainties in $B \to \pi\pi$ see ref. [19].

In conclusion, our analysis shows that measurements of CP asymmetries in several channels may allow the extraction of the CP mixing phase and to disentangle SM and SUSY contributions to the CP decay phase. The golden-plated decays in this respect are $B \to \phi K_S$ and $B \to K_S \pi^0$ channels.

Incl.	Excl.	ϕ_{SM}^D	r_{SM}	ϕ_{SUSY}^D	r_{250}	r_{500}
$b \to c\bar{c}s$	$B \to J/\psi K_S$	0	–	ϕ_{23}	$0.03 - 0.1$	$0.008 - 0.04$
$b \to s\bar{s}s$	$B \to \phi K_S$	0	–	ϕ_{23}	$0.4 - 0.7$	$0.09 - 0.2$
$b \to u\bar{u}s$		P 0				
	$B \to \pi^0 K_S$		$0.01 - 0.08$	ϕ_{23}	$0.4 - 0.7$	$0.09 - 0.2$
$b \to d\bar{d}s$		T γ				
$b \to c\bar{u}d$		0				
	$B \to D_{CP}^0 \pi^0$		0.02	–	–	–
$b \to u\bar{c}d$		γ				
	$B \to D^+ D^-$	T 0	$0.03 - 0.3$		$0.007 - 0.02$	$0.002 - 0.006$
$b \to c\bar{c}d$				ϕ_{13}		
	$B \to J/\psi \pi^0$	P β	$0.04 - 0.3$		$0.007 - 0.03$	$0.002 - 0.008$
	$B \to \phi \pi^0$	P β	–		$0.06 - 0.1$	$0.01 - 0.03$
$b \to s\bar{s}d$				ϕ_{13}		
	$B \to K^0 \bar{K}^0$	u-P γ	$0 - 0.07$		$0.08 - 0.2$	$0.02 - 0.06$
$b \to u\bar{u}d$	$B \to \pi^+ \pi^-$	T γ	$0.09 - 0.9$	ϕ_{13}	$0.02 - 0.8$	$0.005 - 0.2$
$b \to d\bar{d}d$	$B \to \pi^0 \pi^0$	P β	$0.6 - 6$	ϕ_{13}	$0.06 - 0.4$	$0.02 - 0.1$
	$B \to K^+ K^-$	T γ	$0.2 - 0.4$		$0.04 - 0.1$	$0.01 - 0.03$
$b\bar{d} \to q\bar{q}$				ϕ_{13}		
	$B \to D^0 \bar{D}^0$	P β	only β		$0.01 - 0.03$	$0.003 - 0.006$

Table 2: CP phases for B decays. ϕ_{SM}^D denotes the decay phase in the SM; T and P denote Tree and Penguin, respectively; for each channel, when two amplitudes with different weak phases are present, one is given in the first row, the other in the last one and the ratio of the two in the r_{SM} column. ϕ_{SUSY}^D denotes the phase of the SUSY amplitude, and the ratio of the SUSY to SM contributions is given in the r_{250} and r_{500} columns for the corresponding SUSY masses.

The size of the SUSY effects is clearly controlled by the the non-diagonal SUSY mass insertions δ_{ij}, which for illustration we have assumed to have the maximal value compatible with the present experimental limits on B_d^0–$\bar{B}_d^0$ mixing.

5 DM and SUSY: a brief comment

We have strong indications that ordinary matter (baryons) is insufficient to provide the large amount of non-shining matter which has been experimentally proven to exist in galactic halos and at the level of clusters of galaxies [20]. In a sense, this might constitute the "largest" indication of new physics beyond the SM. This statement holds true even after the recent stunning developments in the search for non-shining baryonic objects. In September 1993 the discovery of massive dark objects ("machos") was announced. After five years of intensive analysis it is now clear that in any case machos cannot account for the whole dark matter of the galactic halos.

It was widely expected that some amount of non-shining baryonic matter could exist given that the contribution of luminous baryons to the energy density of the Universe $\Omega = \rho/\rho_{cr}$ ($\rho_{cr} = 3H_0^2/8\pi G$ where G is the gravitational constant and H_0 the Hubble constant) is less than 1%, while from nucleosynthesis we infer $\Omega_{baryon} = \rho_{baryon}/\rho_{cr} = (0.06 \pm 0.02)h_{50}^{-2}$, where $h_{50} = H_0/50$ Km/s Mpc. On the other hand, we have direct indications that Ω should be at least 20% which means that baryons can represent not more than half of the entire energy density of the Universe [20].

We could make these considerations on the insufficiency of the SM to obtain a large enough Ω more dramatic if we accept the theoretical input that the Universe underwent some inflationary era which produced $\Omega =1$. In that case, at least 90% of the whole energy density of the Universe should be provided by some new physics beyond the SM.

Before discussing possible particle physics candidates, it should be kept in mind that DM is not only called for to provide a major contribution to Ω, but also it has to provide a suitable gravitational driving force for the primordial density fluctuations to evolve into the large-scale structures (galaxies, clusters and superclusters of galaxies) that we observe today [20]. Here we encounter the major difficulties when dealing with the two "traditional" sources of DM: Cold (CDM) and Hot (HDM) DM.

Light neutrinos in the eV range are the most typical example of HDM, being their decoupling temperature of O(1 MeV). On the other hand, the Lightest Supersymmetric Particle (LSP) in the tens of GeV range is a typical CDM candidate. Taking the LSP to be the lightest neutralino, one obtains that when it decouples it is already non-relativistic, being its decoupling temperature typically one order of magnitude below its mass.

Both HDM and CDM have some difficulty to correctly reproduce the experimental spectrum related to the distribution of structures at different scales. The conflict is more violent in the case of pure HDM. Neutrinos of few eV's tend to produce too many superlarge structures. The opposite problem arises with pure CDM: we obtain too much power in the spectrum at low mass scales (galactic scales).

A general feature is that some amount of CDM should be present in any case. A possibility which has been envisaged is that after all the whole Ω could be much smaller than one, say 20% or so and then entirely due to CDM. However, if one keeps on demanding the presence of an inflationary epoch, then it seems unnatural to have Ω so different from unity (although lately some variants of inflationary schemes leading to Ω smaller than one have been proposed). Another possibility is that CDM provides its 20% to Ω, while all the rest to reach the unity value is given by a nonvanishing cosmological constant.

Finally, the possibility which encounters quite some interest is the so-called Mixed Dark Matter (MDM) [21], where a wise cocktail of HDM and CDM is present. An obvious realization of a MDM scheme is a variant of the MSSM where neutrinos get a mass of few eV's. In that case the lightest neutralino (which is taken to be the LSP) plays the role of CDM and the light neutrino(s) that of HDM. With an appropriate choice of the parameters it is possible to obtain contributions to Ω from the CDM and HDM in the desired range.

In the MSSM with R parity the lightest SUSY particle (LSP) is absolutely stable. For several reasons the lightest neutralino is the favourite candidate to be the LSP fulfilling the role of CDM [22].

The neutralinos are the eigenvectors of the mass matrix of the four neutral fermions partners of the W_3, B, H_1^0 and H_2^0. There are four parameters entering this matrix: M_1, M_2, μ and $\tan\beta$. The first two parameters denote the coefficients of the SUSY breaking mass terms $\tilde{B}\tilde{B}$ and $\tilde{W}_3\tilde{W}_3$ respectively. μ is the coupling of the $H_1 H_2$ term in the superpotential. Finally

15

$\tan \beta$ denotes the ratio of the VEV's of the H_2 and H_1 scalar fields.

In general M_1 and M_2 are two independent parameters, but if one assumes that grand unification takes place, then at the grand unification scale $M_1 = M_2 = M_3$, where M_3 is the gluino mass at that scale. Then at M_W one obtains:

$$M_1 = \frac{5}{3} \tan^2 \theta_w M_2 \simeq \frac{M_2}{2}, \qquad M_2 = \frac{g_2^2}{g_3^2} m_{\tilde{g}} \simeq \frac{m_{\tilde{g}}}{3}, \qquad (2)$$

where g_2 and g_3 are the SU(2) and SU(3) gauge coupling constants, respectively.

The above relation between M_1 and M_2 reduces to three the number of independent parameters which determine the lightest neutralino composition and mass: $\tan \beta, \mu$ and M_2. Hence, for fixed values of $\tan \beta$ one can study the neutralino spectrum in the (μ, M_2) plane. The major experimental inputs to exclude regions in this plane are the request that the lightest chargino be heavier than M_Z and the limits on the invisible width of the Z hence limiting the possible decays $Z \to \chi\chi, \chi\chi'$.

Let us focus now on the role played by χ as a source of CDM. χ is kept in thermal equilibrium through its electroweak interactions not only for $T > m_\chi$, but even when T is below m_χ. However for $T < m_\chi$ the number of $\chi's$ rapidly decrease because of the appearance of the typical Boltzmann suppression factor $\exp(-m_\chi/T)$. When T is roughly $m_\chi/20$ the number of χ diminished so much that they do not interact any longer, i.e. they decouple. Hence the contribution to Ω_{CDM} of χ is determined by two parameters: m_χ and the temperature at which χ decouples (T_D). T_D fixes the number of $\chi's$ which survive. As for the determination of T_D itself, one has to compute the χ annihilation rate and compare it with the cosmic expansion rate [23].

Several annihilation channels are possible with the exchange of different SUSY or ordinary particles, $\tilde{f}$, H, Z, etc. Obviously the relative importance of the channels depends on the composition of χ. For instance, if χ is a pure gaugino, then the $\tilde{f}$ exchange represents the dominant annihilation mode.

Quantitatively [24], it turns out that if χ results from a large mixing of the gaugino ($\tilde{W}_3$ and $\tilde{B}$) and Higgsino ($\tilde{H}_1^0$ and $\tilde{H}_2^0$) components, then the annihilation is too efficient to allow the surviving χ to provide Ω large enough. Typically in this case $\Omega < 10^{-2}$ and hence χ is not a good CDM candidate. On the contrary, if χ is either almost a pure Higgsino or a pure gaugino then it can give a conspicuous contribution to Ω

In the case χ is mainly a gaugino (say at least at the 90% level) what is decisive to establish the annihilation rate is the mass of $\tilde{f}$. If sfermions are light the χ annihilation rate is fast and the Ω_χ is negligible. On the other hand, if $\tilde{f}$ (and hence $\tilde{l}$, in particular) is heavier than 150 GeV, the annihilation rate of χ is sufficiently suppressed so that Ω_χ can be in the right ballpark for Ω_{CDM}. In fact if all the $\tilde{f}'s$ are heavy, say above 500 GeV and for $m_\chi \ll m_{\tilde{f}}$, then the suppression of the annihilation rate can become even too efficient yielding Ω_χ unacceptably large.

In the minimal SUSY standard model there are five new parameters in addition to those already present in the non–SUSY case. Imposing the electroweak radiative breaking further reduces this number to four. Finally, in simple supergravity realizations the soft parameters A and B are related. Hence we end up with only three new, independent parameters. One can use the constraint that the relic χ abundance provides a correct Ω_{CDM} to restrict the allowed area in this 3–dimensional space. Or, at least, one can eliminate points of this space which would lead to $\Omega_\chi > 1$, hence overclosing the Universe. For χ masses up to 150 GeV it is possible to find sizable regions in the SUSY parameter space where Ω_χ acquires interesting values for the DM problem. A detailed discussion on this point is beyond the scope of this talk. The interested reader can find a thorough analysis in the review of Ref. [22] and the original papers therein quoted.

There exist two ways to search for the existence of relic neutralinos. First we have direct detection: neutralinos interact with matter both through coherent and spin dependent effects. Only coherent effects are currently accessible to direct detection. The sensitivity of the direct detection experiment has reached now an area of the SUSY parameter space of the MSSM which is of great interest for neutralinos in the 50 GeV - 200 GeV range.

The indirect detection is based on the search for signals coming from pair annihilation of neutralinos. Such annihilation may occur inside celestial bodies (Earth, Sun, etc.) where neutralinos may be gravitationally captured. The signal is then a flux of muon neutrinos which can be detected as up-going muons in a neutrino telescope. Another possibility is that the neutralino annihilation occurs in the galactic halo. In this case the signal consists of photon, positron and antiproton fluxes. They can be observed by detectors placed on balloons or satellites. The computation of these fluxes is strongly affected by the composition of the lightest neutralino. In any case also these indirect searches for relic neutralinos are now probing interesting areas of the

MSSM parameter space.

A very different prospect for DM occurs in the GMSB schemes. In this case the gravitino mass ($m_{3/2}$) loses its role of fixing the typical size of soft breaking terms and we expect it to be much smaller than what we have in models with a hidden sector. Indeed, given the well-known relation [1] between $m_{3/2}$ and the scale of SUSY breaking $\sqrt{F}$, i.e. $m_{3/2} = O(F/M)$, where M is the reduced Planck scale, we expect $m_{3/2}$ in the KeV range for a scale $\sqrt{F}$ of $O(10^6$ GeV) that has been proposed in models with low-energy SUSY breaking in a visible sector.

A gravitino of that mass behaves as a Warm Dark Matter (WDM) particle, that is, a particle whose free streaming scale involves a mass comparable to that of a galaxy, $\sim 10^{11-12} M_\odot$.

However, critical density models with pure WDM are known to suffer for serious troubles [25]. Indeed, a WDM scenario behaves much like CDM on scales above λ_{FS}. Therefore, we expect in the light gravitino scenario that the level of cosmological density fluctuations on the scale of galaxy clusters ($\sim 10\,h^{-1}$ Mpc) to be almost the same as in CDM. As a consequence, the resulting number density of galaxy clusters is predicted to be much larger than what observed.

We have recently considered different variants of a light gravitino DM dominated model. It seems that in all cases there exist difficulties to account correctly for cosmic straucture formation. This provides severe cosmological constraints on the GMSB models [26].

In conclusion SUGRA models with R parity offer the best candidate for CDM. It is remarkable that as a by-product of the MSSM we obtain a lightest neutralino which can provide the correct amount of DM in a wide area of the SUSY parametr space. Even more interesting, we are now experimentally approaching the level of sensitivity which is needed to explore (directly or indirectly) large portions of this area of parameter space. The complementarity of this exploration to that performed by using FCNC and CP tests and direct collider SUSY searches looks promising.

Acknowledgements

References

[1] For a phenomenologically oriented review, see:
P. Fayet and S. Ferrara, *Phys. Rep.* **32C** (1977) 249;
H.P. Nilles, *Phys. Rep.* **110** (1984) 1.
H.E. Haber and G.L Kane, *Phys. Rep.* **117** (1987) 1;
For spontaneously broken N=1 supergravity, see:
E. Cremmer, S. Ferrara, L. Girardello and A. Van Proeyen, *Nucl. Phys.* **B 212** (1983) 413 and references therein.
P. Nath, R. Arnowitt and A.H. Chamseddine, *Applied N=1 Supergravity* (World Scientific, Singapore, 1984);
A.G. Lahanas and D.V. Nanopoulos, *Phys. Rep.* **145** (1987) 1.

[2] M. Dine, W. Fischler and M. Srednicki, *Nucl. Phys.* **B 189** (1981) 575;
S. Dimopoulos and S. Raby, *Nucl. Phys.* **B 192** (1981) 353;
M. Dine and W. Fischler, *Phys. Lett.* **B 110** (1982) 227;
M. Dine and M. Srednicki, *Nucl. Phys.* **B 202** (1982) 238;
M. Dine and W. Fischler, *Nucl. Phys.* **B 204** (1982) 346;
L. Alvarez-Gaumé, M. Claudson and M. Wise, *Nucl. Phys.* **B 207** (1982) 96;
C. Nappi and B. Ovrut, *Phys. Lett.* **B 113** (1982) 175;
S. Dimopoulos and S. Raby, *Nucl. Phys.* **B 219** (1983) 479.

[3] A. Nelson and M. Dine, *Phys. Rev.* **D 48** (1993) 1277;
M. Dine, A. E. Nelson and Y. Shirman, *Phys. Rev.* **D 51** (1995) 1362;
M. Dine, A. Nelson, Y. Nir and Y. Shirman, *Phys. Rev.* **D 53** (1996) 2658.

[4] E. Poppitz and S. Trivedi, *Phys. Rev.* **D 55** (1997) 5508;
N. Arkani-Hamed, J. March-Russel and H. Murayama, *Nucl. Phys.* **B 509** (1998) 3;
H. Murayama, *Phys. Rev. Lett.* **79** (1997) 18;
S. Dimopoulos, G. Dvali, G. Giudice and R. Rattazzi, *Nucl. Phys.* **B 510** (1998) 12 ;
S. Dimopoulos, G. Dvali and R. Rattazzi, *Phys. Lett.* **B 413** (1997) 336;
M. Luty, *Phys. Lett.* **B 413** (1997) 71;
T. Hotta, K.-I. Izawa and T. Yanagida, *Phys. Rev.* **D 55** (1997) 415;
N. Haba, N. Maru and T. Matsuoka, *Nucl. Phys.* **B 497** (1997) 31;
L. Randall, *Nucl. Phys.* **B 495** (1997) 37;
Y. Shadmi, *Phys. Lett.* **B 405** (1997) 99;
N. Haba, N. Maru and T. Matsuoka, *Phys. Rev.* **D 56** (1997) 4207;
C. Csaki, L. Randall and W. Skiba, *Phys. Rev.* **D 57** (1998) 383;
Y. Shirman, *Phys. Lett.* **B 417** (1998) 281;
For a complete review, see: G. Giudice and R. Rattazzi, hep-ph/9801271.

[5] R. Barbieri and G. Giudice, *Nucl. Phys.* **B 306** (1988) 63;
G. Anderson and D. Castano, *Phys. Rev.* **D 52** (1995) 1693;
P.H. Chankowski, J. Ellis and S. Pokorski, *Phys. Lett.* **B 423** (1998) 327;
R. Barbieri and A. Strumia, hep-ph/9801353.

[6] H. Haber and R. Hempfling, *Phys. Rev. Lett.* **66** (1991) 1815;
J. Ellis, G. Ridolfi and F. Zwirner, *Phys. Lett.* **B 257** (1991) 83;
Y. Okada, M. Yamaguchi and T. Yanagida, *Prog. Theor. Phys.* **85** (1991) 1.

[7] S. Dimopoulos and S. Thomas, *Nucl. Phys.* **B 465** (1996) 23.

[8] P.H. Chankowski and S. Pokorski, hep-ph/9707497;
D. Pierce and J. Erler, hep-ph/9708374.

[9] J. Ellis and D.V. Nanopoulos, *Phys. Lett.* **B 110** (1982) 44;
R. Barbieri and R. Gatto, *Phys. Lett.* **B 110** (1982) 211.

[10] L.J. Hall, V.A. Kostelecky and S. Raby, *Nucl. Phys.* **B 267** (1986) 415.

[11] M.J. Duncan, *Nucl. Phys.* **B 221** (1983) 285;
J.F. Donoghue, H.P. Nilles and D. Wyler, *Phys. Lett.* **B 128** (1983) 55;
A. Bouquet, J. Kaplan and C.A. Savoy, *Phys. Lett.* **B 148** (1984) 69.

[12] F. Gabbiani and A. Masiero, *Nucl. Phys.* **B 322** (1989) 235;
J.S. Hagelin, S. Kelley and T. Tanaka, *Nucl. Phys.* **B 415** (1994) 293.

[13] E. Gabrielli, A. Masiero and L. Silvestrini, *Phys. Lett.* **B 374** (1996) 80.

[14] F. Gabbiani, E. Gabrielli, A. Masiero and L. Silvestrini, *Nucl. Phys.* **B 477** (1996) 321.

[15] J.A. Bagger, K.T. Matchev and R.-J. Zhang, *Phys. Lett.* **B 412** (1997) 77.

[16] M. Misiak, S. Pokorski and J. Rosiek, hep-ph/9703442.

[17] For some recent discussions tackling this question, see:
N. Deshpande, B. Dutta and S. Oh, *Phys. Rev. Lett.* **77** (1996) 4499;
J. Silva and L. Wolfenstein, *Phys. Rev.* **D 53** (1997) 5331;
A. Cohen, D. Kaplan, F. Leipentre and A. Nelson, *Phys. Rev. Lett.* **78** (1997) 2300;
Y. Grossman and M. Worah, *Phys. Lett.* **B 395** (1997) 241;
Y. Grossman, Y. Nir and R. Rattazzi, hep-ph/9701231;
M. Ciuchini, E. Franco, G. Martinelli, A. Masiero and L. Silvestrini, *Phys. Rev. Lett.* **79** (1997) 978;
Y. Grossman, Y. Nir and M. Worah, *Phys. Lett.* **B 407** (1997) 307;
R. Barbieri and A. Strumia, *Nucl. Phys.* **B 508** (1997) 3.

[18] M. Ciuchini, E. Franco, G. Martinelli, A. Masiero and L. Silvestrini, in ref. [17].

[19] M. Ciuchini, E. Franco, G. Martinelli and L. Silvestrini, *Nucl. Phys.* **B 501** (1997) 271.

[20] For an introduction to the DM problem, see, for instance:
R. Kolb and S. Turner, *The Early universe* (addison-Wesley, New York, N.Y.) 1990;
Dark Matter, ed. by M. Srednicki (North-Holland, Amsterdam) 1989;

J. Primack, D. Seckel and B. Sadoulet, *Ann. Rev. Nucl. Part. Sci.* **38** (1988) 751.

[21] Q. Shafi and F.W. Stecker, *Phys. Lett. B* **53** (1984) 1292;
S.A. Bonometto and R. Valdarnini, *Astroph. J.* **299** (1985) L71;
S. Achilli, F. Occhionero and R. Scaramella, *Astroph. J.* **299** (1985) L77;
J.A. Holtzman, *Astroph. J. Suppl.* **71** (1981) 1;
A.N.Taylor and M. Rowan-Robinson, *Nature* **359** (1992) 396;
J.A. Holtzman and J. Primack, *Astroph. J.* **396** (1992) 113;
D. Pogosyan and A. Starobinski, *Astroph. J.* **447** (1995) 465;
A. Klypin, J. Holtzman, J. Primack and E. Regos, *Astroph. J.* **415** (1993) 1.

[22] G. Jungman, M. Kamionkowski and K. Griest, *Phys. Rep.* **267** (1996) 195, and references therein.

[23] J. Ellis, J.S. Hagelin, D.V. Nanopoulos, K. Olive and M. Srednicki, *Nucl. Phys.* **B 238** (1984) 453.

[24] A. Bottino, F. Donato, N. Fornengo and S. Scopel, hep-ph/9710295.

[25] S. Colombi, S. Dodelson and L.M. Widrow, astro-ph/9505029.

[26] E. Pierpaoli, S. Borgani, A. Masiero and M. Yamaguchi, *Phys. Rev.* **D 57** (1998) 2089.