Antideuterons as a Signature of Supersymmetric Dark Matter

Fiorenza Donato\(^a\) *, Nicolao Fornengo\(^b\) and Pierre Salati\(^{a,c}\)

\(^a\) Laboratoire de Physique Théorique LAPTH, BP110, F-74941 Annecy-le-Vieux Cedex, France.
\(^b\) Instituto de Física Corpuscular - C.S.I.C., Departamento de Física Teórica, Universitat de València, C./ Dr Moliner 50, E-46100 Burjassot, València, Spain.
\(^c\) Université de Savoie, BP1104 73011 Chambéry Cedex, France.

Draft Version of April 29, 1999

Once the energy spectrum of the secondary component is well understood, measurements of the antiproton cosmic–ray flux at the Earth will be a powerful way to indirectly probe for the existence of supersymmetric relics in the galactic halo. Unfortunately, it is still spoilt by considerable theoretical uncertainties. As shown in this work, searches for low–energy antideuterons appear in the mean time as a plausible alternative, worth being explored. Above a few GeV/n, a dozen spallation antideuterons should be collected by the future AMS experiment on board ISSA. For energies less than \(\sim 3\) GeV/n, the \(D\) spallation component becomes negligible and may be supplanted by a potential supersymmetric signal. If a few low–energy antideuterons are discovered, this should be seriously taken as a clue for the existence of massive neutralinos in the Milky Way.

I. INTRODUCTION.

Cosmic ray fluxes are about to be measured with unprecedented precision both by balloon borne detectors and by space instruments. The various ongoing experiments are also hunting for traces of antimatter in the cosmic radiation. The BESS collaboration [1] plans to push the limit on the \(\text{He}/\text{He}\) ratio down to \(10^{-8}\) whereas the AMS spectrometer should reach a sensitivity of \(\sim 10^{-9}\) once it is installed on the International Space Station Alpha (ISSA) [2]. The search for antinuclei has profound cosmological implications. The discovery of a single antihelium or anticarbon would actually be a smoking gun for the existence of antimatter islands in our neighborhood. However, light antinuclei, mostly antiprotons but also antideuterons, are actually produced in our Galaxy as secondaries. They result from the interaction of high–energy cosmic–ray protons with the interstellar gas of the Milky Way disk. In a previous analysis, Chardonnet et al. [3] have estimated the flux of antideuterium \(\bar{D}\) and antihelium \(\bar{3}\)He secondaries. The \(\bar{D}\) signal is very weak but may marginally be detected by AMS on board ISSA. The case of antihelium is, at least for the moment, hopeless.

The dark matter of the Milky Way could be made mostly of elementary particles such as the heavy and neutral species predicted by supersymmetry. The mutual annihilations of these relics, potentially concealed in the halo of our Galaxy, would therefore produce an excess in the cosmic radiation of gamma rays, antiprotons and positrons. In particular, supersymmetric antiprotons should be abundant at low energy, a region where the flux of \(p\) secondaries is a priori negligible. There is quite an excitement trying to extract from the observations a possible \(p\) exotic component which would signal the presence of supersymmetric dark matter in the Galaxy. Unfortunately, it has been recently realized [4–6] that a few processes add up together to flatten out, at low energy, the spectrum of secondary antiprotons. Ionisation losses as well as inelastic but non-annihilating scatterings on the hydrogen atoms of the galactic disk result into the decrease of the antiproton energy. The low–energy tail of the \(\bar{p}\) spectrum is replenished by the more abundant population from higher energies. That effect is further strengthened by solar modulation which also shifts the energy spectrum towards lower energies. As a result of these effects, the secondary \(\bar{p}\)’s are much more abundant at low energy.

---

\hat{\text{E–mail: donato@lapp.in2p3.fr, fornengo@flamenco.ific.uv.es, salati@lapp.in2p3.fr}}\)
than previously thought. Disentangling an exotic supersymmetric contribution from the conventional component of spallation antiprotons may turn out to be a very difficult task. The antiproton signal of supersymmetric dark matter is therefore in jeopardy.

Antideuterons, i.e., the nuclei of antideuterium, are free from such problems. As explained in Sect. II, they form when an antiproton and an antineutron merge together. The two antinucleons must be at rest with respect to each other in order for fusion to take place successfully. For kinematic reasons, a spallation reaction creates very few low–energy particles. Low-energy secondary antideuterons are even further suppressed. Energy loss mechanisms are also less efficient in shifting the antideuteron energy spectrum towards low energies. The corresponding interstellar (IS) flux is derived in Sect. III, for energies in the range extending from 0.1 up to 100 GeV/n. It reaches a maximum of $2 - 5 \times 10^{-8} \text{D m}^{-2} \text{s}^{-1} \text{sr}^{-1} \text{GeV}^{-1}$ for a kinetic energy of $\sim 4 \text{ GeV/n}$. A dozen of secondary antideuterons should be collected by the AMS/ISSA experiment.

On the other hand, supersymmetric D’s are manufactured at rest with respect to the Galaxy. In neutralino annihilations, antinucleons are predominantly produced with low energies. This feature is further enhanced by their subsequent fusion into antideuterons, hence a fairly flat spectrum for supersymmetric antideuterium nuclei as shown in Sect. IV. Below a few GeV/n, secondary antideuterons are quite suppressed with respect to their supersymmetric partners. That low–energy suppression is orders of magnitude more effective for antideuterons than for antiprotons. This makes cosmic–ray antideuterons a much better probe of supersymmetric dark matter than antiprotons.

Unfortunately, antideuteron fluxes are quite small with respect to p’s. We nevertheless show in Sect. V that a significant portion of the supersymmetric parameter space may be explored by measuring the cosmic–ray D flux at low energy. In particular, an AMS/ISSA caliber experiment should reach a sensitivity of $4.8 \times 10^{-8} \text{D m}^{-2} \text{s}^{-1} \text{sr}^{-1} \text{GeV}^{-1}$ at solar minimum, pushing it down to $3.2 \times 10^{-8} \text{D m}^{-2} \text{s}^{-1} \text{sr}^{-1} \text{GeV}^{-1}$ at solar maximum, for a modulated energy of 0.24 GeV/n.

II. PRODUCTION OF ANTIDEUTERONS.

At this point, our goal is to derive the cross section for the production of antideuterons. The processes at stake are both the spallation of a cosmic–ray high–energy proton on an hydrogen atom at rest and the annihilation of a neutralino pair. The number $dN_X$ of particles X – antinucleons or antideuterons – produced in a single reaction and whose momenta are $\vec{k}_X$, is related to the differential production cross section through

$$dN_X = \frac{1}{\sigma_{\text{tot}}} d^3 \sigma_X(\sqrt{s}, \vec{k}_X),$$

where $\sigma_{\text{tot}}$ denotes the total cross section for the process under scrutiny – spallation reaction or neutralino annihilation. The total available energy is $\sqrt{s}$. The corresponding differential probability for the production of X is defined as

$$dN_X = \mathcal{F}_X(\sqrt{s}, \vec{k}_X) d^3 k_X.$$

(2)

For each of the processes under concern, the differential probability for the production of an antiproton or an antineutron may be derived. The calculation of the probability for the formation of an antideuteron can now proceed in two steps. We first need to estimate the probability for the creation of an antiproton–antineutron pair. Then, those antinucleons merge together to yield an antinucleus of deuterium.

As explained in Ref. [3], the production of two antinucleons is assumed to be proportionnal to the square of the momenta of the two particles. The hypothesis that factorization of the probabilities holds is fairly well established at high energies. For spallation reactions, however, the bulk of the antiproton production takes place for an energy $\sqrt{s} \sim 10$ GeV which turns out to be of the same order of magnitude as the antideuteron mass. Pure factorization should break in that case as a result of energy conservation. It needs to be slightly adjusted. We have therefore assumed that the center of mass energy available for the production of the second antinucleon is reduced by twice the energy carried away by the first antinucleon

$$\mathcal{F}_{\bar{p}\bar{n}}(\sqrt{s}, \vec{k}_{\bar{p}}, \vec{k}_{\bar{n}}) = \frac{1}{2} \mathcal{F}_{\bar{p}}(\sqrt{s}, \vec{k}_{\bar{p}}) \mathcal{F}_{\bar{n}}(\sqrt{s} - 2E_{\bar{p}}, \vec{k}_{\bar{n}}) + (\vec{k}_{\bar{p}} \leftrightarrow \vec{k}_{\bar{n}}).$$

(3)

Once the antiproton and the antineutron are formed, they combine together to give an antideuteron with probability

$$\mathcal{F}_{\bar{D}}(\sqrt{s}, \vec{k}_{\bar{D}}) d^3 k_{\bar{D}} = \int d^3 k_{\bar{p}} d^3 k_{\bar{n}} c(k_{\bar{p}}, k_{\bar{n}}) \mathcal{F}_{\bar{p}\bar{n}}(\sqrt{s}, \vec{k}_{\bar{p}}, \vec{k}_{\bar{n}}).$$

(4)
The summation is performed on those antinucleon configurations for which

\[ k_\tilde{p} + k_\tilde{n} = k_\tilde{D}. \]  

The coalescence function \( C(k_\tilde{p}, k_\tilde{n}) \) describes the probability for a \( \tilde{p} - \tilde{n} \) pair to yield an antideuteron. That function depends actually on the difference \( k_\tilde{p} - k_\tilde{n} = 2\Delta \) between the antinucleon momenta so that relation (4) may be expressed as

\[ \mathcal{F}_\tilde{D}(\sqrt{s}, k_\tilde{D}) = \int d^3\Delta C(\Delta) \mathcal{F}_{\tilde{p},\tilde{n}} \left( \sqrt{s}, k_\tilde{p} = \frac{k_\tilde{D}}{2} + \Delta, k_\tilde{n} = \frac{k_\tilde{D}}{2} - \Delta \right). \]  

An energy of \( \sim 3.7 \text{ GeV} \) is required to form by spallation an antideuteron whereas the binding energy of the latter is \( B \sim 2.2 \text{ MeV} \). The coalescence function is therefore strongly peaked around \( \Delta = 0 \) and expression (6) simplifies into

\[ \mathcal{F}_\tilde{D}(\sqrt{s}, k_\tilde{D}) \approx \left\{ \int d^3\Delta C(\Delta) \right\} \mathcal{F}_{\tilde{p},\tilde{n}} \left( \sqrt{s}, k_\tilde{p} = \frac{k_\tilde{D}}{2}, k_\tilde{n} = \frac{k_\tilde{D}}{2} \right), \]  

where the probability for the formation of the \( \tilde{p} - \tilde{n} \) pair has been factored out. The term in brackets may be estimated in the rest frame of the antideuteron through the Lorentz invariant term

\[ \int \frac{E_\tilde{D}}{E_\tilde{p} E_\tilde{n}} d^3\Delta C(\Delta) \approx \left( \frac{m_\tilde{D}}{m_\tilde{p} m_\tilde{n}} \right) \left( \frac{4}{3} \pi P_{\text{coal}}^3 \right). \]  

In that frame, the antinucleons merge together if the momentum of the corresponding two–body reduced system is less than some critical value \( P_{\text{coal}} \). That coalescence momentum is the only free parameter of our factorization and coalescence scheme. As shown in Ref. [3], the resulting antideuteron production cross section in proton–proton collisions is well fitted by this simple one–parameter model. A value of \( P_{\text{coal}} = 58 \text{ MeV} \) has been derived, not too far from what may be naively expected from the antideuteron binding energy, \( \sqrt{m_\tilde{p} B} \sim 46 \text{ MeV} \).

The differential probability with which an antiproton is produced during a proton–proton collision is related to the corresponding Lorentz invariant cross section through

\[ \sigma_{\tilde{p} \rightarrow \tilde{n}}^{\text{tot}} |_{E_\tilde{p}} \times F_{\tilde{n}} \left( \sqrt{s}, k_\tilde{p} \right) = E_\tilde{p} \frac{d^3\sigma}{d^3k_\tilde{p}} |_{LI}. \]  

The latter is experimentally well known. It is fairly well fitted by the Tan and Ng’s parametrization [7] which has been used here. Assuming that the invariance of isospin holds, the antineutron production cross section is equal to its antiproton counterpart. The Lorentz invariant cross section for the production of antideuterons resulting from the impact of a high–energy cosmic–ray proton on a proton at rest has been derived by Chardonnet et al. [3] who showed that

\[ E_\tilde{D} \frac{d^3\sigma_\tilde{D}}{d^3k_\tilde{D}} = \left( \frac{m_\tilde{D}}{m_\tilde{p} m_\tilde{n}} \right) \left( \frac{4}{3} \pi P_{\text{coal}}^3 \right) \times \frac{1}{2\sigma_{\tilde{p} \rightarrow \tilde{n}}} \times \]  

\[ \times \left\{ E_\tilde{p} \frac{d^3\sigma_\tilde{p}}{d^3k_\tilde{p}} \left( \sqrt{s}, k_\tilde{p} \right) E_n \frac{d^3\sigma_\tilde{n}}{d^3k_\tilde{n}} \left( \sqrt{s} - 2E_\tilde{p}, k_\tilde{n} \right) + \left( k_\tilde{p} \leftrightarrow k_\tilde{n} \right) \right\}. \]  

The corresponding differential cross section obtains from the summation, in the galactic frame, of the Lorentz invariant production cross section (10)

\[ \frac{d\sigma_{\text{ph–}D}}{dE_\tilde{D}} \left( E_p \rightarrow E_\tilde{D} \right) = 2\pi k_\tilde{D} \int_{0}^{\theta_{\text{max}}} E_\tilde{D} \frac{d^3\sigma}{d^3k_\tilde{D}} |_{LI} d(-\cos \theta). \]  

In that frame, \( \theta \) denotes the angle between the momenta of the incident proton and of the produced antideuteron. It is integrated up to a maximal value \( \theta_{\text{max}} \) set by the requirement that, in the center of mass frame of the reaction, the antideuteron energy \( E_{\tilde{D}} \) cannot exceed the bound
\[ E_{\overline{D}}^{\text{max}} = \frac{s - 16m_p^2 + m_D^2}{2\sqrt{s}}. \]  

(12)

The integral (11) is performed at fixed antideuteron energy \( E_{\overline{D}}^2 = m_D^2 + k_{\overline{D}}^2 \).

In the case of a neutralino annihilation, the differential multiplicity for antiproton production may be expressed as

\[
\frac{dN_{\overline{p}}}{dE_{\overline{p}}} = \sum_{F,h} B_{\chi h}^{(F)} \frac{dN_{\overline{h}}}{dE_{\overline{p}}}.
\]

(13)

The annihilation proceeds, through the various final states \( F \), towards the quark or the gluon \( h \) with the branching ratio \( B_{\chi h}^{(F)} \). Quarks or gluons may be directly produced when a neutralino pair annihilates. They may alternatively result from the intermediate production of a Higgs or gauge boson as well as of a top quark. Each quark or gluon \( h \) generates in turn a jet whose subsequent fragmentation and hadronization yields the antiproton energy spectrum \( dN_{\overline{h}}/dE_{\overline{p}} \). Because neutralinos are at rest with respect to each other, the probability to form, say, an antiproton with momentum \( k_{\overline{p}} \) is essentially isotropic

\[
\frac{dN_{\overline{p}}}{dE_{\overline{p}}} (\chi + \chi \rightarrow \overline{p} + \ldots) = 4\pi k_{\overline{p}} E_{\overline{p}} f_p(\sqrt{s} = 2m_\chi, E_{\overline{p}}).
\]

(14)

Applying the factorization–coalescence scheme discussed above leads to the antideuteron differential multiplicity

\[
\frac{dN_{\overline{D}}}{dE_{\overline{D}}} = \left( \frac{4P_{\text{coal}}}{3k_{\overline{D}}} \right) \left( \frac{m_D}{m_\overline{p} m_{\overline{n}}} \right) \sum_{F,h} B_{\chi h}^{(F)} \left( \frac{dN_{\overline{h}}}{dE_{\overline{p}}} \left( E_{\overline{p}} = E_{\overline{D}}/2 \right) \right)^2.
\]

(15)

It may be expressed as a sum, extending over the various quarks and gluons \( h \) as well as over the different annihilation channels \( F \), of the square of the antiproton differential multiplicity. That sum is weighted by the relevant branching ratios. The antineutron and antiproton differential distributions have been assumed to be identical. The hypothesis that factorization holds is certainly conservative. The antinucleons which merge together to create an antideuteron are produced in the same quark or gluon jet. Their momenta are not isotropically distributed with respect to each other. They tend to be more aligned than what has been assumed here, with a larger chance to generate an antideuteron. However, our analysis is meant to be conservative.

III. THE DETECTION OF SPALLATION ANTIDEUTERONS.

As suggested by Parker, the propagation of cosmic–rays inside the Galaxy is strongly affected by their scattering on the irregularities of magnetic fields. This results into a diffusive transport. In the following, we will assume an isotropic diffusion with an empirical value for the diffusion coefficient. Our Galaxy can be reasonably well modelled by a thin disk of atomic and molecular hydrogen, with radius \( R \sim 20 \) kpc and thickness \( \sim 200 \) pc. This gaseous ridge is sandwiched between two diffusion regions which act as confinement domains as a result of the presence of irregular magnetic fields. They extend vertically up to \( \sim 3 \) kpc apart from the central disk. That two-zone diffusion model is in good agreement with the observed primary and secondary nuclei abundances [8].

Assuming a steady regime, the propagation of cosmic–ray antideuterons within the Milky Way is accounted for by the diffusion equation

\[
-K \Delta \psi_{\overline{D}} + \Gamma_{\overline{D}} \psi_{\overline{D}} + \frac{\partial}{\partial E} \left\{ b(E) \psi_{\overline{D}} \right\} = q_{\text{sec}}^{\overline{D}},
\]

(16)

where \( \psi_{\overline{D}} \) is the density of antideuterons per unit of volume and per unit of energy. In the left–hand side of relation (16), the first term describes the diffusion of the particles throughout the galactic magnetic fields. The coefficient \( K \) is derived from measurements of the light element abundances in cosmic–rays. It is constant at low energies, but beyond a critical value of \( R_0 = 1 \) GV, it raises with rigidity \( R \) like

\[
K(R) = K_0 \left( 1 + \frac{R}{R_0} \right)^{0.6},
\]

(17)
where $K_0 = 6 \times 10^{27}$ cm$^2$ s$^{-1}$. It is assumed to be essentially independent of the nature of the species that propagate throughout the Galaxy. The second term accounts for the destruction of antideuterons through their interactions, mostly annihilations, with the interstellar medium. Antideuterons may also undergo fragmentation if they survive annihilation. In that case, they are broken apart as most of the cosmic-ray nuclei. The total collision rate is given by

$$\Gamma_{\bar{D}} = \sigma_{\bar{D}H} v_{\bar{D}} n_H,$$

(18)

where $\sigma_{\bar{D}H}$ is the total antideuteron interaction cross section with protons $[9]$, $v_{\bar{D}}$ denotes the velocity and $n_H = 1$ cm$^{-3}$ is the average hydrogen density in the thin matter disk. The last term in the left-hand side of relation (16) stands for the energy losses undergone by antideuterons as they diffuse in the galactic ridge. The rate $b(E_{\bar{D}})$ at which the antideuteron energy varies is essentially set by the ionization losses which the particle undergoes as it travels through interstellar gas. This mechanism yields the following contribution to the energy loss rate

$$b_{\text{ion}}(E_{\bar{D}}) = -4\pi r_e^2 m_e c^2 n_H \frac{e}{\beta} \left\{ \ln \left( \frac{2m_e c^2}{E_0} \right) + \ln \left( \beta^2 \gamma^2 \right) - \beta^2 \right\}.$$

(19)

In molecular hydrogen, the ionization energy $E_0$ has been set equal to 19.2 eV; here $\gamma = E_{\bar{D}}/m_{\bar{D}}$. The classical radius of the electron is denoted by $r_e$ and the electron mass is $m_e$. In the case of antiprotons, it was realized $[7,4,5]$ that the dominant energy loss mechanism is actually their inelastic, but non-annihilating, interactions with interstellar protons. The latter are excited towards resonant states and hence absorb part of the antiproton energy. In the $p$ frame, an incident proton kicks off the antiproton at rest, transferring some of its kinetic energy. In the case of antideuterons, however, such a process is no longer possible. In the $\bar{D}$ frame, the impinging proton cannot

![Graph showing the IS secondary flux of antideuterons as a function of kinetic energy per nucleon. The solid curve corresponds to the median value of the cosmic-ray proton spectrum, as derived by Bottino et al. The dashed and dotted lines respectively stand for the maximal and minimal values of the primary proton flux from which the antideuterons originate.](image)
FIG. 2. The median IS spectrum of Fig. 1 (solid curve) has been modulated at solar maximum (dashed line) and minimum (dotted line).

transfer energy without destroying the antideuteron whose binding energy \( B \sim 2.2 \text{ MeV} \) is much smaller than the typical kinetic energies at stake. That is why fragmentation generally dominates the interactions of cosmic–ray nuclei with interstellar matter. Accordingly, the resulting destruction occurs at fixed energy per nucleon. In the right–hand side of the diffusion Eq. (16), the production rate \( q_{\text{sec}}^{\bar{D}} \) of the spallation antideuterons involves a convolution over the incident cosmic–ray proton energy spectrum \( \psi_p \) of the differential production cross section (11)

\[
q_{\text{disk}}^{\bar{D}}(E_{\bar{D}}) = \int_{E_{\bar{D}}}^{+\infty} dE_p \psi_p(E_p) n_H v_p \frac{d\sigma_{pH \rightarrow \bar{D}}}{dE_{\bar{D}}} \left\{ E_p \rightarrow E_{\bar{D}} \right\} . \tag{20}
\]

The differential energy distribution \( \psi_{\bar{D}} \) of secondary antideuterons is determined by solving Eq. (16). We have followed the standard approach which may be found in Ref. [10]. At the edge of the domain where the cosmic–rays are confined, the particles escape freely, the diffusion becomes inefficient and densities vanish. This provides the boundary conditions for solving Eq. (16). Then, because the problem is axisymmetric, the various cosmic–ray distributions may be expanded as series of Bessel functions of zeroth order. Details may be found in Refs. [4,11]. The secondary antideuteron interstellar flux finally obtains from the differential energy spectrum

\[
\Phi_{\text{sec}}^{\bar{D}} = \frac{1}{4\pi} \psi_{\bar{D}} v_{\bar{D}} . \tag{21}
\]

The interstellar (IS) flux of spallation antideuterons is presented in Fig. 1 as a function of the kinetic energy per nucleon. As explained in Bottino et al. [4], the IS proton flux is still uncertain around \( \sim 20–100 \text{ GeV} \), an energy range that contributes most to the integral (20). We have borrowed the parametrization

\[
\Phi_{p}^{\text{IS}} = \beta \left( \frac{E_p}{1 \text{ GeV}} \right)^{-\alpha} . \tag{22}
\]
The median IS proton flux corresponds to a normalization factor of $A = 15,950$ protons m$^{-2}$ s$^{-1}$ sr$^{-1}$ GeV$^{-1}$ with a spectral index of $\alpha = 2.76$. The normalization factor $A$ has been varied from $12,300$ (minimal) up to $19,600$ protons m$^{-2}$ s$^{-1}$ sr$^{-1}$ GeV$^{-1}$ (maximal). Accordingly, the minimal and maximal IS proton fluxes respectively correspond to the spectral indices $\alpha = 2.61$ and $2.89$. In Fig. 1, the solid curve features the IS secondary antideuterons generated from the median proton spectrum. The maximal (dashed line) and minimal (dotted line) distributions delineate the band within which the spallation antideuteron signal lies. The flux reaches a maximum value comprised between $2.1$ and $4.9 \times 10^{-8}$ D m$^{-2}$ s$^{-1}$ sr$^{-1}$ GeV$^{-1}$ for a kinetic energy of $\sim 4$ GeV/n. The antideuteron spectrum sharply drops below a few GeV/n. Remember that in the galactic frame, the production threshold is $17 m_p$. When a high–energy cosmic–ray proton impinges on an hydrogen atom at rest, the bulk of the resulting antiprotons and antineutrons keep moving, with kinetic energies $\sim 10 - 20$ GeV. For kinematical reasons, the production of antinucleons at rest with respect to the Galaxy is extremely improbable. The manufacture of a low–energy antideuteron is even more improbable. It actually requires the creation of both an antiproton and an antineutron at rest. The momenta need to be aligned in order for fusion to successfully take place. Low–energy antideuterons produced as secondaries in the collisions of high–energy cosmic–rays with the interstellar material are therefore extremely scarce, with a completely depleted energy spectrum below $\sim 1$ GeV/n. Energy losses tend to shift the antideuteron spectrum towards lower energies with the effect of replenishing the low–energy tail with the more abundant species which, initially, had a higher energy. This process tends to slightly soften the strong decrease of the low–energy antideuteron spectrum. The effect is nevertheless mild. Remember that in the case of antiprotons, it is actually the inelastic but non–annihilating interactions which considerably flatten the $\bar{p}$ distribution. Ionizations losses are not enough to significantly affect the energy spectrum. The IS secondary antideuterons are therefore extremely depleted below $\sim 1$ GeV/n. The spallation background is negligible in the region where supersymmetric $\bar{D}$'s are expected to be most abundant. This feature makes the detection of low–energy antideuterons an interesting signature of the presence of supersymmetric relics in the Galaxy.

In Fig. 2, the median IS $\bar{D}$ spectrum (solid curve) has been modulated at solar maximum (dashed line) and solar minimum (dotted line). We have applied the forced field approximation [12] to estimate the effect of the solar wind on the cosmic–ray energies and fluxes. For the energies at stake, this amounts to simply shift the IS energy of a nucleus $N$, with charge $Z$ and atomic number $A$, by a factor of $Z\Phi$. The solar modulation parameter $\Phi$ has the same dimensions as a rigidity or an electric potential. The Earth ($\oplus$) and interstellar (IS) energies, per nucleon, are therefore related by

$$E_N^\oplus/A = E_N^{\text{IS}}/A - |Z|e\Phi/A \ .$$

(23)

In Perko’s approximation, antinuclei are affected in just the same way as nuclei. Their energy decreases as they penetrate inside the heliomagnetic field. Once the momenta at the Earth $p_N^\oplus$ and at the boundaries of the heliosphere $p_N^{\text{IS}}$ are determined, the flux modulation ensues

$$\Phi_N^\oplus(E_N^\oplus)/\Phi_N^{\text{IS}}(E_N^{\text{IS}}) = \left\{\frac{p_N^\oplus}{p_N^{\text{IS}}}\right\}^2 \ .$$

(24)

Antideuterons undergo an energy loss, per nucleon, half that of protons and antiprotons. At solar minimum (maximum) the modulation parameter $\Phi$ has been set equal to $320$ MV ($800$ MV) [4]. The energy shift is larger at solar maximum than at solar minimum. Once modulated, the sharply decreasing IS antideuteron distribution tends therefore to be flatter at solar maximum as is clear in Fig. 2. We estimate that a total of 12–13 secondary antideuterons may be collected by the AMS collaboration during the space station stage, in the energy range extending up to 100 GeV/n. These antideuterons correspond to IS energies in excess of $\sim 3$ GeV/n, a region free from the effects of solar modulation. This result takes into account the geomagnetic suppression as discussed in Sect. V.

IV. THE SUPERSYMMETRIC ANTIDEUTERON SIGNAL.

As a theoretical framework, we use the Minimal Supersymmetric extension of the Standard Model (MSSM) [13], which conveniently describes the supersymmetric phenomenology at the electroweak scale, without too strong theoretical assumptions. This model has been largely adopted by many authors for evaluations of the neutralino relic abundance and detection rates (for reviews, see Refs. [14,15]).

The MSSM is defined at the electroweak scale as a straightforward supersymmetric extension of the Standard Model. The Higgs sector consists of two Higgs doublets $H_1$ and $H_2$ and, at the tree level, is fully described by two free
FIG. 3. The IS flux of secondary antideuterons (heavier solid curve) decreases at low energy whereas the energy spectrum of the antideuterons from supersymmetric origin tends to flatten. The four cases of table I are respectively featured by the solid (a), dotted (b), dashed (c) and dot-dashed (d) curves.

TABLE I. These four cases illustrate the richness of the supersymmetric parameter space. There is no obvious correlation between the antiproton and antideuteron Earth fluxes with the neutralino mass $m_{\chi}$. Case (c) is a gaugino-higgsino mixture and still yields signals comparable to those of case (a), yet a pure gaugino. Antideuteron fluxes are estimated at both solar minimum and maximum, for a modulated energy of 0.24 GeV/n. The last column features the corresponding number of $\bar{D}$'s which AMS on board ISSA can collect below 3 GeV/n.

<table>
<thead>
<tr>
<th>case</th>
<th>$m_\chi$</th>
<th>$P_\beta$ (%)</th>
<th>$\Omega_{\chi}h^2$</th>
<th>$\Phi_{\min}^{\bar{D}}$ (0.24 GeV/n)</th>
<th>$\Phi_{\max}^{\bar{D}}$ (0.24 GeV/n)</th>
<th>$\Phi_{\max}^{\bar{D}}$ (0.24 GeV/n)</th>
</tr>
</thead>
<tbody>
<tr>
<td>$a$</td>
<td>36.5</td>
<td>96.9</td>
<td>0.20</td>
<td>$1.2 \times 10^{-3}$</td>
<td>$1.0 \times 10^{-7}$</td>
<td>$2.9 \times 10^{-8}$</td>
</tr>
<tr>
<td>$b$</td>
<td>61.2</td>
<td>95.3</td>
<td>0.13</td>
<td>$3.9 \times 10^{-3}$</td>
<td>$3.5 \times 10^{-7}$</td>
<td>$1.1 \times 10^{-7}$</td>
</tr>
<tr>
<td>$c$</td>
<td>90.4</td>
<td>53.7</td>
<td>0.03</td>
<td>$1.1 \times 10^{-3}$</td>
<td>$1.8 \times 10^{-7}$</td>
<td>$6.1 \times 10^{-8}$</td>
</tr>
<tr>
<td>$d$</td>
<td>120</td>
<td>98.9</td>
<td>0.53</td>
<td>$2.9 \times 10^{-4}$</td>
<td>$2.5 \times 10^{-8}$</td>
<td>$8.6 \times 10^{-9}$</td>
</tr>
</tbody>
</table>

parameters, namely: the ratio of the two vacuum expectation values $\tan \beta \equiv \langle H_2 \rangle / \langle H_1 \rangle$ and the mass of one of the three neutral physical Higgs fields, which we choose to be the mass $m_A$ of the neutral pseudoscalar one. Once radiative corrections are introduced, the Higgs sector depends also on the squark masses through loop diagrams. The radiative corrections to the neutral and charged Higgs bosons, adopted in the present paper, are taken from Refs. [16,17]. The other parameters of the model are defined in the superpotential, which contains all the Yukawa interactions and the Higgs–mixing term $\mu H_1 H_2$, and in the soft–breaking Lagrangian, which contains the trilinear and bilinear breaking parameters and the soft gaugino and scalar mass terms. In this model, the neutralino is defined as the lowest–mass linear superposition of photino ($\tilde{\gamma}$), zino ($\tilde{Z}$) and the two higgsino states ($\tilde{H}_1^0$, $\tilde{H}_2^0$).
\[ \chi \equiv a_1 \bar{\gamma} + a_2 \bar{Z} + a_3 \bar{H}_1^0 + a_4 \bar{H}_2^0 . \]  

In order to deal with manageable models, it is necessary to introduce some assumptions which establish relations among the too many free parameters at the electroweak scale. We adopt the following usual conditions. All trilinear parameters are set to zero except those of the third family, which are unified to a common value \( A \). All squarks and sleptons soft–mass parameters are taken as degenerate: \( m_{\tilde{l}} = m_{\tilde{q}} = m_0 \). The gaugino masses are assumed to unify at \( M_{\text{GUT}} \), and this implies that the \( U(1) \) and \( SU(2) \) gaugino masses are related at the electroweak scale by \( M_1 = (5/3) \tan^2 \theta_W M_2 \). When all these conditions are imposed, the supersymmetric parameter space is completely described by six independent parameters, which we choose to be: \( M_2, \mu, \tan \beta, m_A, m_0, A \). In our analyses, we vary them in the following ranges: 20 GeV \( \leq M_2 \leq 500 \) GeV; 20 GeV \( \leq |\mu| \leq 500 \) GeV; 80 GeV \( \leq m_A \leq 1000 \) GeV; 100 GeV \( \leq m_0 \leq 1000 \) GeV; \( -3 \leq A \leq +3 \); 1 \( \leq \tan \beta \leq 50 \).

The supersymmetric parameter space is constrained by all the experimental limits achieved at accelerators on supersymmetric and Higgs searches [18]. Also the constraints due to the \( b \rightarrow s + \gamma \) process [19,20] have been taken into account (see Ref. [21] for a discussion of our implementation of the \( b \rightarrow s + \gamma \) constraint and for the relevant references). We further require the neutralino to be the Lightest Supersymmetric Particle (LSP) and the supersymmetric configurations to provide a neutralino relic abundance in accordance with the cosmological bound \( \Omega_\chi h^2 \leq 0.7 \) [14].

For the evaluation of the averaged annihilation cross section \(< \sigma_{\text{ann}} v >\), we have followed the procedure outlined in Ref. [22]. We have considered all the tree–level diagrams which are responsible for neutralino annihilation and which are relevant to \( \bar{p} \) production, namely: annihilation into quark–antiquark pairs, into gauge bosons, into a Higgs boson pair and into a Higgs and a gauge boson. For each final state we have considered all the relevant Feynman diagrams, which involve the exchange of Higgs and \( Z \) bosons in the s–channel and the exchange of squarks, neutralinos and charginos in the t and u–channels. Finally, we have included the one–loop diagrams which produce a two–gluon final state [23]. The \( \bar{p} \) differential distribution \( dN_{\bar{p}}/dE_{\bar{p}} \) has been evaluated as discussed in Ref. [22]. Here we only recall that we have calculated the branching ratios \( B_{\chi h}^{(F)} \) for all annihilation final states \( F \) which may produce \( \bar{p} \)'s. These final states fall into two categories: (i) direct production of quarks and gluons and (ii) generation of quarks through the intermediate production of Higgs bosons, gauge bosons and \( t \) quark. In order to obtain the distributions \( dN_{\bar{p}}/dE_{\bar{p}} \), the hadronization of quarks and gluons has been computed by using the Monte Carlo code Jetset 7.2 [24]. For the top quark, we have considered it to decay before hadronization. The source term for supersymmetric antideuterons

\[ dN_{\bar{D}}^\text{asy} (\chi + \chi \rightarrow \bar{D} + \ldots) = < \sigma_{\text{ann}} v > \frac{dN_{\bar{D}}}{dE_{\bar{D}}} \left( \frac{\rho_{\chi}}{m_{\chi}} \right)^2 \]  

(26)
FIG. 5. The supersymmetric–to–secondary IS flux ratio for antiprotons (lower curves) and antideuterons (upper curves) is presented as a function of the kinetic energy per nucleon. The supersymmetric configurations are those reported in Table I and featured in Figs. 3 and 4. Below a few GeV/n, the flux ratio is always larger for \( \bar{D} \)'s than for \( \bar{p} \)'s. For the supersymmetric configurations of Table I, the antiproton signal is swamped into its background. This is not the case for antideuterons. At low energy, the flux of primaries is several orders of magnitude above the \( \bar{D} \) background.

Supplements the spallation contribution \( q_{\bar{D}}^{sec} \) in the diffusion Eq. (16). The propagation of primary antideuterons from the remote regions of the galactic halo to the Earth has been treated as explained in Ref. [4]. The neutralino distribution has been assumed to be spherical, with radial dependence

\[
\rho_\chi = \rho_\chi^\odot \left\{ \frac{a^2 + r_\odot^2}{a^2 + m^2} \right\},
\]

where \( m^2 = r^2 + z^2 \). The solar system is at a distance \( r_\odot \) of 8 kpc from the galactic center. The dark matter halo has a core radius \( a = 3.5 \) kpc and its density in the solar neighborhood is \( \rho_\chi^\odot = 0.4 \) GeV cm\(^{-3} \) [14].

In Fig. 3, both primary (supersymmetric) and secondary (spallation) interstellar antideuterons energy spectra are presented. The secondary flux (heavier solid line) drops sharply at low energies as discussed above. The four supersymmetric examples of Table I are respectively featured by the solid (a), dotted (b), dashed (c) and dot-dashed (d) curves. The corresponding primary fluxes flatten at low energy where they reach a maximum. As the secondary \( \bar{D} \) background vanishes, the supersymmetric signal is the largest. Neutralino annihilations actually take place at rest in the galactic frame. The fragmentation and subsequent hadronization of the jets at stake tend to favour the production of low–energy species. Therefore, the spectrum of supersymmetric antiprotons – and antineutrons – is fairly flat below \( \sim 1 \) GeV. For the same reasons, the coalescence of the primary antideuterons produced in neutralino annihilations predominantly takes place with the two antinucleons at rest, hence a flat spectrum at low energy, as is clear in Fig. 3. The fusion of an antideuteron requires actually that its antinucleon constituents should be aligned in momentum space. Consequently, secondary antideuterons are completely depleted below \( \sim 1 \) GeV while the primary species are mostly produced in that low–energy regime. This trend still appears once the energies and fluxes are
modulated. The left and right panels of Fig. 4 respectively show the effects of solar modulation at maximum and minimum. The spallation background somewhat flattens. It is still orders of magnitude below the supersymmetric signal which clearly exhibits a plateau.

It is difficult to establish a correlation between the \( \hat{\mathbb{D}} \) flux and the neutralino mass. In case (c), for instance, \( m_\chi \) is \( \sim 3 \) times larger as in case (a) and yet the corresponding antideuteron flux is larger. It is not obvious either that gaugino–like mixtures lead to the largest \( \hat{\mathbb{D}} \) signals. Table I gives a flavor of the complexity and of the richness of the supersymmetric parameter space.

In Fig. 5, the supersymmetric–to–spallation IS flux ratio for antiprotons (lower curves) and antideuterons (upper curves) are presented as a function of the kinetic energy per nucleon. In the case of antiprotons, the primary–to–secondary ratio is much smaller than for antideuterons. For the configurations of table I presented here, the \( \hat{\mathbb{p}} \) primary flux is at the same level as the spallation background. The supersymmetric antiproton signal is swamped in the flux of the secondaries. This is not the case for antideuterons. At low energies, their supersymmetric flux is several orders of magnitude above background. Antideuterons appear therefore as a much cleaner probe of the presence of supersymmetric relics in the galactic halo than antiprotons. The price to pay however is a much smaller flux. Typical \( \hat{\mathbb{D}} \) spectra may reach up to \( 10^{-6} - 10^{-5} \text{ m}^{-2} \text{ s}^{-1} \text{ sr}^{-1} \text{ GeV}^{-1} \). This corresponds to an antiproton signal of \( 10^{-2} - 10^{-1} \text{ m}^{-2} \text{ s}^{-1} \text{ sr}^{-1} \text{ GeV}^{-1} \), \( i.e., \) four orders of magnitude larger. It is therefore crucial to ascertain which portion of the supersymmetric configurations will be accessible to future experiments through the detection of low–energy cosmic–ray antideuterons.

V. DISCUSSION AND CONCLUSIONS.

In order to be specific, we have estimated the amount of antideuterons which may be collected by the AMS experiment once it is on board ISSA. The future space station is scheduled to orbit at 400 km above sea level, with an inclination of \( \alpha = 52^\circ \) with respect to the Earth equator. A revolution takes about 1.5 hours so that ISSA should fly over the same spot every day. The AMS detector may be pictured as a cylindrical magnetic field with diameter \( D = 110 \text{ cm} \). At any time, its axis points towards the local vertical direction. The colatitude of the north magnetic pole has been set equal to \( \Upsilon = 11^\circ \). At any given time \( t \) along the orbit, the geomagnetic latitude \( \varrho \) of ISSA may be inferred from

\[
\sin \varrho = \sin \Upsilon \cos (\Omega_{\text{sid}} t) \cos (\Omega_{\text{orb}} t + \varphi) + \\
+ \cos \alpha \sin \Upsilon \sin (\Omega_{\text{sid}} t) \sin (\Omega_{\text{orb}} t + \varphi) + \sin \alpha \cos \Upsilon \sin (\Omega_{\text{orb}} t + \varphi)
\]

where \( \Omega_{\text{sid}} \) and \( \Omega_{\text{orb}} \) respectively denote the angular velocities associated to the sideral rotation of the Earth and to the orbital motion of the space station. The phase \( \varphi \) depends on the orbital initial conditions and does not affect the result if a large number of revolutions – typically 100 – is considered. The Earth is shielded from cosmic–rays because its magnetic field prevents particles from penetrating downwards. At any given geomagnetic latitude \( \varrho \), there exists a rigidity cut–off \( R_{\text{min}} \) below which the cosmic–ray flux is suppressed. This lower bound depends on the radius \( R \) of the orbit through

\[
R_{\text{min}} = \frac{\mu_\oplus \cos^4 \varrho}{\varpi^2},
\]

where \( \mu_\oplus \) denotes the Earth magnetic dipole moment and \( \mu_\oplus / R_\oplus^2 \simeq 60 \text{ GV} \). The term \( \varpi \) stands for

\[
\varpi = 1 + \sqrt{1 + \cos \theta \cos^3 \varrho}.
\]

It depends on the angle \( \theta \) between the cosmic–ray momentum at the detector and the local east–west line that points in the ortho–radial direction of an axisymmetric coordinate system. Notice that because we are interested here in singly charged species, the rigidity amounts to the momentum \( p \). Once the cosmic–ray energy as well as the geomagnetic latitude are specified, the solid angle \( \Omega_{\text{cut}} \) inside which the direction of the incoming particle lies may be derived from relations (29) and (30). The AMS detector looks upwards within \( \sim 27^\circ \) around the vertical. This corresponds to a solid angle of \( \Omega_{\text{det}} = 0.68 \text{ sr} \). Because the apparatus does not point towards the local east or west, impinging particles may not be seen by the instrument. The effective solid angle \( \Omega_{\text{eff}} \) through which they are potentially detectable corresponds to the overlap, if any, between \( \Omega_{\text{cut}} \) and \( \Omega_{\text{det}} \). The value of \( \Omega_{\text{eff}} \) depends on the cosmic–ray rigidity \( p \) as well as on the precise location of the detector along the orbit. The detector acceptance may therefore be defined as
$\mathcal{R}(p) = \frac{\pi}{4} D^2 \int \Omega_{\text{eff}}(p,t) \, dt$, \hspace{1cm} (31)

where the time integral runs over the duration $\tau$ of the space mission. In the case of AMS on board ISSA, $\tau$ is estimated to be $10^8$ s (3 yrs). Between 100 MeV/n and 100 GeV/n, we infer a total acceptance of $5.8 \times 10^9$ m$^2$ s sr GeV for antiprotons and of $6 \times 10^9$ m$^2$ s sr GeV for antideuterons. The net number of cosmic-ray species which AMS may collect on board ISSA is actually a convolution of the detector acceptance with the relevant differential flux at Earth. For antideuterons, this leads to

$$N_D = \int \mathcal{R}(p^\oplus D) \Phi_D(dT^\oplus_D),$$ \hspace{1cm} (32)

where the integral runs on the $D$ modulated energy $T^\oplus_D$.

Integrating the secondary flux discussed in Sect. III leads respectively to a total of 12.3 and 13.4 antideuterons, depending on whether the solar cycle is at maximum or minimum. These spallation $D$'s are mostly expected at high energies. As is clear from Figs. 3 and 4, the secondary flux drops below the supersymmetric signal below a few GeV/n. The transition typically takes place for an interstellar energy of 3 GeV/n. Below that value, the secondary antideuteron signal amounts to a total of only 0.6 (solar maximum) and 0.8 (solar minimum) nuclei. Most of the supersymmetric signal is therefore concentrated in a low–energy band extending from the AMS threshold of 100 MeV/n up to a modulated energy of 2.6 GeV/n (maximum) or 2.84 GeV/n (minimum) which corresponds to an upper bound of 3 GeV/n in interstellar space. In this low–energy region where spallation antideuterons yield a negligible background, the AMS acceptance is $2.2 \times 10^7$ m$^2$ s sr GeV for antiprotons and $5.5 \times 10^7$ m$^2$ s sr GeV for antideuterons.

FIG. 6. The supersymmetric $\bar{D}$ flux has been integrated over the range of IS energies extending from 0.1 up to 3 GeV/n. The resulting yield $N^\oplus_D$ of antideuterons which AMS on board ISSA can collect is plotted as a function of the neutralino mass $m_\chi$. Modulation has been considered at solar maximum.
For each supersymmetric configuration, the $\bar{D}$ flux has been integrated over that low–energy range. The resulting yield $N_{\bar{D}}$ which AMS may collect on board ISSA is presented as a function of the neutralino mass $m_\chi$ in the scatter plot of Fig. 6. During the AMS mission, the solar cycle will be at maximum. Most of the configurations are gaugino like (crosses) or mixed combinations of gaugino and higgsino states (dots). A significant portion of the parameter space is associated to a signal exceeding one antideuteron – horizontal dashed line. In a few cases, AMS may even collect more than a dozen of low–energy $\bar{D}$ nuclei. However, when the antideuteron signal exceeds $\sim 20$ particles, the associated antiproton flux is larger than what BESS 95 + 97 [26] has measured.

The scatter plot of Fig. 6 may be translated into a limit on the antideuteron flux $\Phi_{\bar{D}}^\oplus$ at the Earth. Table I gives a flavor of the relation between that flux and the yield $N_{\bar{D}}$ of low–energy antideuterons. At solar maximum, a value of $N_{\bar{D}} = 1$ translates, on average, into a flux of $\sim 3.2 \times 10^{-8}$ $\bar{D}$ m$^{-2}$ s$^{-1}$ sr$^{-1}$ GeV$^{-1}$ for a modulated energy of 240 MeV/n. The energy spectrum matters of course. For the steep differential flux of case (a), a value of $4.8 \times 10^{-8}$ $\bar{D}$ m$^{-2}$ s$^{-1}$ sr$^{-1}$ GeV$^{-1}$ is necessary in order to achieve a signal of at least one antideuteron. In case (d) where the spectrum is much flatter, the same $\bar{D}$ yield is reached for a flux of only $2.8 \times 10^{-8}$ $\bar{D}$ m$^{-2}$ s$^{-1}$ sr$^{-1}$ GeV$^{-1}$. The horizontal dashed lines of Figs. 7 should therefore be understood as averaged limits. They are nevertheless indicative of the level of sensitivity which may be reached through the search for low–energy antideuterons. The left and right panels respectively correspond to a solar activity taken at maximum and minimum. In these scatter plots, the $\bar{D}$ modulated flux is featured as a function of the neutralino mass $m_\chi$. The antideuteron energy at the Earth has been set equal to 240 MeV/n. The flux $\Phi_{\bar{D}}^\oplus$ is larger at solar minimum – when modulation is weaker – than at maximum. The lower the cosmic–ray energy, the larger that effect. The plateaux of Figs. 4 illustrate the flatness of the supersymmetric $\bar{D}$ spectra at low energies. These plateaux actually exhibit a shift by a factor $\sim 3$ between the left and right panels. Accordingly, the constellation of supersymmetric antideuterons collected at low energy obtains from the convolution of Eq. (32). It also varies during the solar cycle, in a somewhat lesser extent however than the above mentioned plateaux. Between maximum and minimum, the value of $N_{\bar{D}}$ only varies by a factor of $\sim 2$, to be compared to a flux increase of $\sim 3$. At solar maximum, when AMS/ISSA will be operating, a signal of one antideuteron translates into a flux sensitivity of $\sim 3.2 \times 10^{-9}$ antinuclei m$^{-2}$ s$^{-1}$ sr$^{-1}$ GeV$^{-1}$. At minimum, the same signal would translate into the weaker limit of $\sim 4.8 \times 10^{-9}$ antideuterons m$^{-2}$ s$^{-1}$ sr$^{-1}$ GeV$^{-1}$ and the horizontal dashed line is shifted upwards by $\sim 50%$. The
supersymmetric configurations which an antideuteron search may unravel are nevertheless more numerous at solar minimum. Between the left and the right panels, the constellation of representative points is actually shifted upwards and, relative to the limit of sensitivity, the increase amounts to a factor $\sim 2$. In spite of the low fluxes at stake, the antideuteron channel is sensitive to a respectable number of supersymmetric configurations.

Supersymmetric antiprotons are four orders of magnitude more abundant in cosmic rays than antideuterons – see table I. However, as already discussed, they may be swamped in the background arising from the secondaries. The AMS experiment will collect a large number of antiprotons on board ISSA. Our concern is whether an hypothetical supersymmetric $\bar{p}$ signal may be disentangled from the background. Because the latter still suffers from large theoretical uncertainties, we are afraid that antiproton searches in cosmic rays are not yet the ultimate probe for the existence of supersymmetric relics in the Milky Way. As discussed in Refs. [4–6], the distribution of secondary antiprotons turns out to be flatter than previously estimated. Therefore, it is still a quite difficult task to ascertain which fraction of the measured antiproton spectrum may be interpreted as a supersymmetric component. Notice however that as soon as the secondary $\bar{p}$ flux is reliably estimated, low-energy antiproton searches will become a more efficient tool. Meanwhile, we must content ourselves with using observations as a mere indication of what a supersymmetric component cannot exceed. The vertical shaded band of Figs. 8 and 9 corresponds actually to the 1–σ antiproton flux which the BESS 95 + 97 experiments [26] have measured at a $\bar{p}$ energy of 0.24 GeV. In Fig. 8, the supersymmetric antideuteron yield $N_{\bar{D}}$ has been derived at solar maximum. This corresponds to the conditions of the future AMS mission on board the space station. The antideuteron yield is plotted as a function of the associated supersymmetric $\bar{p}$ flux at Earth. The latter is estimated at solar minimum to conform to the BESS data to which the vertical band refers. The scatter plot of Fig. 8 illustrates the strong correlation between the antideuteron and antiproton signals, as may be directly guessed from Eq. 15. The horizontal dashed line indicates the level of

![Graph showing antideuteron yield ($N_{\bar{D}}$) as a function of supersymmetric antiproton flux ($\Phi_{\bar{p}}$) at solar maximum.](image)
sensitivity which AMS/ISSA may reach. Points located above that line but on the left of the shaded vertical band are supersymmetric configurations that are not yet excluded by antiproton searches and for which the antideuteron yield is potentially detectable. The existence of such configurations illustrates the relevance of an antideuteron search at low energies. As shown in Fig. 9, the number of interesting configurations is largest at solar minimum. Both $D$ and $\bar{p}$ fluxes at Earth are plotted against each other. Energies have been set equal to 0.24 GeV/n. The correlation between the antideuteron and antiproton cosmic–ray fluxes is once again noticeable.

Once the energy spectrum of the secondary component is no longer spoilt by considerable theoretical uncertainties, measurements of the antiproton cosmic–ray flux will be a powerful way to test the existence of supersymmetric relics in the galactic halo. In the mean time, searches for low–energy antideuterons appear as a plausible alternative, worth being explored. A dozen spallation antideuterons should be detected by the future AMS experiment on board ISSA above a few GeV/n. For energies less than $\sim$ 3 GeV/n, the $D$ spallation component becomes negligible and may be supplanted by a potential supersymmetric signal. We conclude that the discovery of a few low–energy antideuterons should be taken seriously as a clue for the existence of massive neutralinos in the Milky Way.

Acknowledgements

We would like to express our gratitude toward S. Bottino for stimulating discussions. We also wish to thank R. Battiston and J.P. Vialle for supplying us with useful information pertaining to the AMS experiment. This work was supported by DGICYT under grant number PB95–1077 and by the TMR network grant ERBFMRXCT960090 of the European Union.