MIXED DARK MATTER AND THE FATE OF BARYON
AND LEPTON SYMMETRIES *

A. Masiero
Istituto Nazionale di Fisica Nucleare
35131 Padova (Italy)

and

Dipartimento di Fisica - Univ. di Perugia
06100 Perugia (Italy)

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Abstract. The available data on large scale structures seem to favour models with mixed dark matter (MDM), i.e. with a hot and cold component in a rather well-defined amount, or with some form of “warm” dark matter. I discuss some prospects for these new scenarios for DM in the context of supersymmetric extensions of the electroweak standard model. In particular, I emphasize the intriguing link which exists between the present prospects of solution of the DM puzzle and the explicit or spontaneous breaking of baryon and/or lepton number symmetries. Some consequences on the issue of baryogenesis are worked out.

1. INTRODUCTION

All the three branches which merge together into the relatively recent field of astroparticle physics exhibit a standard model. In particle physics this is the extraordinarily successful Glashow–Weiberg–Salam description of electroweak interactions, in astrophysics we have the standard picture of stellar evolution and in cosmology the hot Big Bang model represents our standard view of the early Universe. Obviously all these three models have become standard thanks to numerous and solid experimental pieces of evidence. In particular, let me remind some particle physicists who consider cosmology rather far from being experimentally testable, that the expansion of the universe, the prediction of the cosmic microwave background radiation and its temperature, and the prediction of the abundances of the primordial elements from nucleosynthesis represent solid experimental pillars which severely constrain any attempt to propose a cosmological model. In fact also in cosmology we are now entering a new phase of observational activity both with large (10-m) ground telescopes and satellite activity.

Given that the three above standard models have common areas, one can naturally wonder whether there is a full compatibility. There is a good chance that this is not the case: the solar neutrino problem may represent the clash between the standard solar model and the standard GWS model (where neutrinos are strictly massless), while the dark matter problem may be the hint of a severe clash between one of the most stringent predictions of nucleosynthesis, the number of surviving baryons, and the absence in the GWS model of relic particles which may be needed to account for the presence of dark matter.

Indeed, I would say that at the moment the electroweak standard model does not feel any serious threat from the accelerator data (the potential discrepancies concerning \( \Gamma(Z \to b\bar{b}) \)), the SLAC data on the left–right asymmetry and the semileptonic branching ratio of the B meson may turn out to be real problems
for the GWS model, but it is certainly premature to draw any conclusion so far. In a sense, the solar neutrino and dark matter problems may represent the only “observational” hint for the need of new physics beyond this model.

In this talk I would like first to discuss to what extent the dark matter problem actually calls for new physics (sect. II). Then, I’ll turn to analyze which kind of new physics may be more suitable for the solution of the DM problem and the related issue of large scale structure formation (sect. III). Sect. IV will be devoted to a study of the relation between DM and the most attractive extension of the SM, i.e. the minimal supersymmetric standard model (MSSM). In Sect. V I’ll deal with the lightest supersymmetric particle as a favourite candidate for cold DM In Sect. VI I’ll introduce a new subject which I think is intimately related to the issue of DM, namely the violation of baryon (B) and lepton (L) numbers at finite temperature and its implications for models with explicit or spontaneous breaking of L. Here the link between two major issues of modern cosmology, i.e. baryogenesis and DM, should appear in all its evidence. In particular this study may be relevant in constraining the different options which are present in the supersymmetrization of the SM. The most recent options for the solution of the DM puzzle with the presence of mixed DM or warm DM will be briefly discussed in Sect. VII.

2. DOES THE DM PROBLEM CALL FOR NEW PHYSICS?

I think that the most relevant question for a particle physicist when tackling the problem of DM is whether the solution of this puzzle calls for extensions of the electroweak SM. Let us briefly state the “facts” [1].

The contribution of luminous matter 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%. The most solid piece of evidence that we need DM comes from the rotation curves of spiral galaxies, with a value of $\Omega_{DM}$ in the 10% range. Given that in the SM the only candidates to produce all this enormous amount of non–shining matter are baryons, one can ask: can we account for $\Omega_{DM}$ just using non–shining baryons? Here comes a crucial constraint on $\Omega_B$ from the Big Bang nucleosynthesis. The ratio of the baryon to photon number densities is one of the three key–elements which established the moment of start of nucleosynthesis and, hence, the abundances of the primordial elements which are produced throughout this process. From detailed analyses one concludes that $\Omega_B$ cannot exceed 10%.
On the other hand there are indications for larger values of $\Omega$ when one applies the usual dynamical methods to scale structures at distances larger than the galactic scale. Last, but certainly not least for a theorist, $\Omega = 1$ is predicted in inflationary models and I think that it is completely fair to say that so far we have not any other viable way to tackle formidable cosmological problems such as causality, oldness and flatness of the Universe, but the inflationary path.

Clearly then, if one believes that $\Omega$ should exceed 10% there is no way at all to accommodate this value of $\Omega$ just involving the presence of the surviving cosmic baryon asymmetry.

In the SM the relics of the primordial Universe are the photons (the famous cosmic background radiation at $2.7 \, ^0\text{K}$), the massless neutrinos (with a number density slightly smaller than that of the photons) and the surviving baryons and charged leptons. Since $\Omega_\gamma$ and $\Omega_\nu$ are certainly much smaller than 10%, we conclude that we need to extend the SM to schemes with additional relic particles if we are to explain $\Omega > 0.1$.

3. WHAT DM IS MADE OF

The first broad distinction among the several candidates for DM which have been proposed in the literature concerns the amount of interactions of a particle with all the others in the primeval plasma. The typical scale one has to compare these interactions with is the expansion rate of the Universe which vastly changes with time. Hence at a certain moment throughout the history of the Universe a particle can exhibit interactions whose rates are larger than the expansion rate of the Universe, while at other times the opposite situation can occur. The former case refers to a situation in which the particle is said to be in thermal equilibrium, whilst in the opposite case we have a particle which is decoupled.

There are particles whose interactions are so weak that they were never in thermal equilibrium. The most representative of these non-thermal candidates is the axion. In this talk I’ll focus my attention on thermal candidates, i.e. particles which were in thermal equilibrium for some time during the early story of the Universe.

The traditional distinction one makes is between hot (HDM) and cold (CDM) dark matter. Two examples can immediately clarify this distinction.

Consider a massive neutrino of few eV’s. The weak interactions keep it in thermal equilibrium as long as the temperature of the Universe is above 1 MeV. Below 1 MeV the neutrino decouples. Hence, at the moment it decouples this neu-
trino is highly relativistic. This is the “standard” example of an HDM candidate. Now, let us envisage a kind of opposite situation. Consider a supersymmetric (SUSY) extension of the SM where the lightest SUSY particle is a neutralino of, say, 50 GeV. As we’ll see in next sect., this particle decouples when the temperature of the Universe is much below its mass (roughly \( \sim m_\chi/20 \), where \( m_\chi \) denotes the mass of the lightest neutralino). Hence at the moment it decouples, this particle is highly non–relativistic. We have here an example of CDM. A more appropriate definition of CDM and HDM is linked to the problem of large scale structure formation which is the subject to which I turn now.

To be a good candidate for DM it is not enough to provide \( \Omega = 1 \), or whichever value of \( \Omega \) one prefers. Very severe constraints on the nature of DM come from the crucial issue of the formation of large scale structures (galaxies, clusters and superclusters of galaxies, etc.). The theory of structure formation is linked to two key–elements: i) the shape of the primordial density fluctuations whose evolution produces the large scale structure that we observe today and ii) the content of matter in the Universe, i.e. the nature of the DM. The variation of these two ingredients leads to different predictions of the power spectrum, i.e. on the distribution of structures at different distances.

Two types of origin for the seed of density fluctuations have been envisaged: inflation and topological defects (cosmic strings,...). In the inflationary scenarios quantum fluctuations of the inflation field are changed into density fluctuations giving rise to a typical scale–invariant fluctuations spectrum. The seed density fluctuations evolve under the action of gravity. Hence their evolution is determined by the nature of DM.

Two scales of importance for the evolution of the seed density fluctuations are: \( \lambda_{FS} \), the free streaming scale below which fluctuations in a nearly collisionless component are damped due to free streaming and \( \lambda_{EQ} \), the horizon length when radiation–matter equality occurs (this scale is important since density fluctuations of non–relativistic matter within the horizon are suppressed during the radiation dominated era, while they begin as the matter domination era starts).

Let us see how how our prototypes for HDM and CDM, the light massive neutrino and the lightest neutralinos, behave in the process of formation of large scale structures.

First I consider light \((m_\nu < 1 \text{ MeV})\) stable neutrinos. If they have a mass \( > 10^{-4} \text{ eV} \) they are non–relativistic today and their energy density is simply given by \( \rho_\nu = m_\nu n_\nu \), where \( m_\nu \) denotes their mass, while \( n_\nu \) is their number.
density. This latter quantity can be easily related to the photon number density $n_\gamma, n_\nu = (3/22)g_\nu n_\gamma$, where $g_\nu$ is equal to 2 or 4 according to the Majorana or Dirac nature of neutrinos. Then one can readily compute the contribution to $\Omega$ due to the presence of these relic neutrinos:

$$\Omega_\nu \equiv \frac{\rho_\nu}{\rho_c} = 0.01 \ m_\nu (eV) \ h_0^{-2} \left(\frac{g_\nu}{2}\right) \left(\frac{T_0}{2.7}\right)^3,$$

(1)

where $h_0$ is the present value of the Hubble parameter in units of $100 \ Km \ sec^{-1} \ parsec^{-1}$ and $T_0$ is the temperature of the microwave cosmic background radiation in degrees Kelvin.

The lower bound of $10^9$ years on the age of the Universe requires $\Omega h_0^2 < 1$ and, therefore:

$$m_\nu (eV) < 200 \ g_\nu^{-1} \ eV$$

(2)

for stable neutrinos which decouple while still relativistic (i.e. $m_\nu < 1 \ MeV$).

Given that experimentally $h_0$ ranges between 0.4 and 1, it is easy to see from (1) that neutrinos in the 10 eV range can readily yield $\Omega_\nu$ in the interesting range $0.1–1$. From this point of view, clearly massive neutrinos would be the best candidates for DM providing large values of $\Omega$ quite easily and with a major advantage on all other competitors: of all the proposed DM candidates, neutrinos are the only particles that we know to exist for sure!

However, as I said, it is not enough to provide $\Omega = 1$ for a relic particle to prove to be a good DM candidate. The other test concerns the role it plays in structures formation. The O (10 eV) neutrinos we are considering are relativistic until late in the evolution of the Universe. The $\nu$ density perturbations are wiped out below the free-streaming scale

$$\lambda_{FS}^\nu \simeq 40 \ Mpc \left(\frac{30 eV}{m_\nu}\right)$$

(3)

corresponding to the mass scale:

$$m_{FS}^\nu \simeq 10^{15} M_\odot \left(\frac{30 eV}{m_\nu}\right),$$

(4)

where $M_\odot$ denotes the solar mass. Hence the first structures to form have dimension much larger than that of galaxies and there is the problem to form enough “small” structures in a scenario with only neutrinos constituting the DM. The only solution which may be viable is the addition to neutrinos of some seeds for
the formation of small structures. Cosmic strings are the best known candidates to play such a role. Whether schemes with pure HDM and cosmic strings may reproduce correctly the known power spectrum is a highly debated issue and the improvement of the current numerical simulations will hopefully shed some light on this intriguing question.

The difficulties which are present in any scheme with pure HDM to account for the structure formation made scenarios with pure CDM even more favoured for several years. The so-called standard cold dark matter model \[^2\] predicted \(\Omega = 1\), with \(\Omega_{CDM} \sim 90 - 95\%\), \(\Omega_B \sim 5 - 10\%\) and \(\Omega_{\nu,\gamma} < 1\%\). The seed fluctuations were generated during inflation and with a scale–invariant spectrum. In this model \(\lambda_{EQ} \simeq 30(\Omega h_0^2)^{-1}\) Mpc. Although some problems were present even before the advent of the COBE data \[^3\], the situation has become rather difficult for the pure CDM scenario after COBE. With the normalization fixed at the COBE data \[^4\] the CDM model predicts more power at small scales than observed \[^5\].

Several remedies have been proposed modifying either the initial fluctuation spectrum or the composition of DM. To “disfavour” the formation of structures at small scales one could try to increase the above value of \(\lambda_{EQ}\). Late decaying particles \[^6\] or a conspicuous contribution of the cosmological constant to \(\Omega\) (with \(\Omega_{CDM} \sim 0.2\) \[^7\]) can yield such an increase of \(\lambda_{EQ}\). The other option to solve the problem is obvious from our previous analysis of the virtues and faults of HDM and CDM scenarios. Since they suffer from opposite problems when dealing with the structure formation, one might expect that a convenient admixture of both components may reproduce the whole power spectrum correctly. It turns out that the best fit is provided by the \(\Omega_{CDM} \sim 0.6\) and \(\Omega_{HDM} \sim 0.3\) \[^8\]. There has been some work along the lines of these mixed dark matter scenarios and some aspects will be discussed in sect. VII. The other possibility that one can envisage is to have a DM candidate which is somewhat “colder” than the above mentioned light neutrinos so that \(\lambda_{FS}\) can decrease. Also some example of this kind of warm DM will be provided in sect. VII.

From the above discussion it emerges that at least some amount of \(\Omega_{DM}\) should be accounted for by the presence of cold dark matter. Before the impressive results of LEP a popular candidate for CDM was a heavy neutrino with a mass in the GeV range. Indeed one can find that \(\Omega_{\nu}h_0^2 \sim 3(\text{GeV}/m_\nu)^2\) and, hence, having \(m_\nu \sim \text{few GeV}\) one could easily obtain \(\Omega_{\nu} \simeq 1\). However if these new heavy neutrinos couple to the \(Z\) boson in the same way ordinary neutrinos do, they would contribute too much to the \(Z\) invisible width. The only way to drastically
reduce this contribution is if these neutrinos have masses close to $m_Z/2$, but in this case $\Omega_\nu$ drops down to 0(1%) making these neutrinos uninteresting for the DM problem.

The favoured CDM candidate has to do with what I consider the most “plausible” extension of the SM, i.e. its supersymmetrization. This is the issue that I intend to discuss in the next sect.

4. DARK MATTER AND SUPERSYMMETRY

There are several reasons which favour the presence of supersymmetry (SUSY) among the fundamental symmetries [9]. In my view the most compelling one is related to the incorporation of gravity with the other three fundamental interactions through supergravity. However, for that matter supersymmetry might as well be a good symmetry at the Planck scale being broken below that scale. If that is the case, then we should not bother so much about SUSY from the phenomenological point of view. What is actually crucial for the TeV physics to be tested in the coming machines is that supersymmetry has to be present much below the Planck scale, indeed down to the electroweak scale of $0(10^2 - 10^3 \text{ GeV})$, if we are to invoke supersymmetry to alleviate the gauge hierarchy problem. As is well known, this problem is related to the presence of fundamental scalar particles in the SM. The most radical cure for the problem would be the elimination altogether of elementary scalars, but then one has to envisage some kind of dynamical mechanism for the spontaneous breaking of the electroweak symmetry. Since so far no consistent model of this kind has been proposed (in spite of years of relentless efforts along these lines), low energy supersymmetry (meaning SUSY extensions of SM with SUSY broken only at $10^2 - 10^3 \text{ GeV}$) represents the only consistent way we have at the moment to cope with the gauge hierarchy problem.

A point of utmost relevance which is often forgotten when discussing the supersymmetrization of the SM is that there is no unique way to realize a SUSY version of the SM. The simplest thing one can try is to use just the fields of the SM embedding them into the convenient superfields and then impose the $SU(3) \times SU(2) \times U(1) \times \text{SUSY}$ symmetry. If one just follows this kind of “minimal prescription”, the model which results is going to be immediately ruled out for a very good reason: your protons would have already decayed before you end reading this sentence! Indeed, one can construct renormalizable operators which violate either baryon (B) or lepton (L) number in the part of the SUSY lagrangian which is known as the superpotential. The latter constitutes a kind of SUSY version
of the ordinary Yukawa lagrangian of the SM, but with a major difference: since in the SUSY version there exist scalar SUSY partners which carry \( B \) or \( L \) it is possible to construct operator of dimension 4 containing two ordinary fermions and one \( s \)-fermion which respect the \( SU(3) \times SU(2) \times U(1) \) symmetry. For instance \( u_R d_R \tilde{d}_R \) and \( u_L e_L \tilde{d}_L \) violate \( B \) and \( L \), respectively (\( \tilde{d}_R \) and \( \tilde{d}_L \) denote the scalar partner of the right–handed down quark or, equivalently, of the left–handed quark \( Q = +1/3 \) down anti–quark). Their simultaneous presence leads to a proton decay through a 4–quark operator mediated by the exchange of a down \( s \)-quark. Since SUSY is bound to be broken at a scale which cannot significantly exceed 1 TeV, we would have an essentially immediate proton decay.

The simplest possibility to avoid the above catastrophe is the addition to the \( SU(3) \times SU(2) \times U(1) \times N = 1 \) SUSY invariance of a new discrete symmetry which forbids all the \( B \) and \( L \) violating operators of the superpotential. This is the famous discrete matter \( R \)–parity which assigns +1 to all known particles of SM and -1 to their SUSY partners. Obviously, then, no operator with two ordinary fermions and one \( s \)-fermion can survive.

This situation that we encounter when supersymmetrizing the SM is profoundly different from what occurs in the SM itself. In this model \( B \) and \( L \) are automatic symmetries of the theory, namely given the \( SU(3) \times SU(2) \times U(1) \) invariance and the usual field assignment it is impossible to construct renormalizable operators which violate \( B \) or \( L \).

\( R \)–parity eliminates all operators which violate \( B \) or \( L \). However, to prevent proton decay it is enough to forbid either \( B \) or \( L \) violation. Hence, one might wonder whether \( R \)–symmetry can be replaced by other discrete symmetries which forbid either the \( B \)– or the \( L \)–violating renormalizable operators, but not all of them. An exhaustive search for all these symmetries was accomplished in ref. [10].

If one imposes the stringent constraint that the \( Z_n \) discrete symmetries which accomplish the task to stop proton decay be “discrete anomalous free”, then one is left with only two candidates: the well–known \( R \)–symmetry and baryon–parity, a discrete symmetry which forbids the \( B \) violating operators, but allows for the \( L \) violating ones. I’ll discuss some aspects of \( B \)–parity in relation to the DM problem in next section.

There is a major implication for the DM issue if one imposes the \( R \)–parity: as long as this symmetry is unbroken the lightest SUSY particle (LSP) is absolutely stable. One can expect that together with \( \gamma, \nu \) and baryons also the LSP will be part of the relics of the early Univers in SUSY versions of the SM with \( R \)–parity.
5. THE LIGHTTEST SUPERSYMMETRIC PARTICLE (LSP)

In models where a discrete symmetry, matter \( R \)-parity \([9]\) discriminates between ordinary and SUSY particles, 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 \([11, 12]\).

The neutralinos are the eigenvectors of the mass matrix of the four neutral fermions partners of the \( W_3, B, H^0 \) and \( H^0_2 \). There are four parameters entering this matrix: \( M_1, M_2, \mu \) and \( tg \beta \). The first two parameters denote the coefficient of the SUSY breaking mass terms \( \bar{B}B \) and \( \bar{W}_3W_3 \) respectively, \( \mu \) is the coupling of the \( H_1 - H_2 \) term the superpotential. Finally \( tg \beta \) 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 a grand unification scale takes place, then at the grand unification \( M_1 = M_2 = M_3 \), where \( M_3 \) is the gluino mass at that scale. Then at \( M_w \) one obtains:

\[
M_1 = \frac{5}{3} tg^2 \theta_w M_2 \simeq \frac{M_2}{2}, \quad M_2 = \frac{g_2^2}{g_3^3} m_{\tilde{g}} \simeq m_{\tilde{g}}/3 ,
\]

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

The relation (5) between \( M_1 \) and \( M_2 \) reduces to three the number of independent parameters which determine the lightest neutralino composition and mass: \( tg \beta, \mu \) and \( M_2 \). Hence, for fixed values of \( tg \beta \) one can study the neutralino spectrum in the \((\mu, M_2)\) plane. The major experimental inputs to exclude regions in this plane are the request that the lightest chargino be heavier than \( M_Z/2 \) and the limits on the invisible width of the \( Z \) hence limiting the possible decays \( Z \to \chi\chi, \chi' \).

Moreover if the GUT assumption is made, then the relation between \( M_2 \) and \( m_{\tilde{g}} \) implies a severe bound on \( M_2 \) from the experimental lower bound on \( m_{\tilde{g}} \) of CDF (roughly \( m_{\tilde{g}} > 120 \) GeV, hence implying \( M_2 > 40 \) GeV). The theoretical demand that the electroweak symmetry be broken radiatively, i.e. due to the renormalization effects on the Higgs masses when going from the superlarge scale of supergravity breaking down to \( M_W \), further constrains the available \((\mu, M_2)\) region.

The first important outcome of this is that the lightest neutralino mass exhibits a lower bound of roughly 10 to 20 GeV \([13]\). The prospects for an improvement of this lower limit at LEP 200 crucially depends on the composition of \( \chi \) \([13]\). If \( \chi \) is mainly a gaugino, then it is difficult to go beyond 40 GeV for such
a lower bound, whilst with a $\chi$ mainly higgsino the lower bound can jump up to $m_\chi > M_W$ at LEP 200.

Let us focus now on the role played by $\chi$ as a source of CDM. $\chi$ is kept in thermal equilibrium through its electroweak interactions not only for $T > m_\chi$, but even when $T$ is below $m_\chi$. 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 $\chi$ diminished so much that they do not interact any longer, i.e. they decouple. Hence the contribution to $\Omega_{CDM}$ of $\chi$ is determined by two parameters: $m_\chi$ and the temperature at which $\chi$ 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 $\chi$ annihilation rate and compare it with the cosmic expansion rate $^{[11]}$.

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 $\chi$. For instance, having assumed $\chi$ to be a pure gaugino in the case discussed in the previous section, then the $\tilde{f}$ exchange represents the dominant annihilation mode.

Quantitatively $^{[14]}$, it turns out that if $\chi$ 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 $\chi$ to provide $\Omega$ large enough. Typically in this case $\Omega < 10^{-2}$ and hence $\chi$ is not a good CDM candidate. On the contrary, if $\chi$ is either almost a pure higgsino or a pure gaugino then it can give a conspicuous contribution to $\Omega$.

As I already mentioned in the previous section, in the case $\chi$ mainly a gaugino (say at least at the 90% level), what is decisive to establish the annihilation rate is the mass of $\tilde{f}$. LEP 200 will be able, hopefully, to test slepton masses up to $M_W$. If there exists a $\tilde{l}$ with mass $< M_W$ then the $\chi$ annihilation rate is fast and the $\Omega_\chi$ is negligible. On the other hand, if $\tilde{f}$ (and hence $\tilde{l}$, in particular) is heavier than 150 GeV, the annihilation rate of $\chi$ is sufficiently suppressed so that $\Omega_\chi$ can be in the right ballpark for $\Omega_{CDM}$. In fact if all the $\tilde{f}$'s are heavy, say above 500 GeV and for $m_\chi << m_{\tilde{f}}$, then the suppression of the annihilation rate can become even too efficient yielding $\Omega_\chi$ unacceptably large. In conclusion if a slepton is found at LEP 200, then the $\chi$ pure gaugino is excluded as a candidate for CDM. If $m_{\tilde{f}}$ is in the range 150 GeV to 500 GeV for $\chi$ in the 20 to 100 GeV range it is possible to give rise to an acceptable value of $\Omega_{CDM}$.

Let us briefly discuss the case of $\chi$ being mainly a higgsino. If the lightest
neutralino is to be predominantly a combination of $\tilde{H}_1^0$ and $\tilde{H}_2^0$ it means that $M_1$ and $M_2$ have to be much larger than $\mu$. Invoking the relation (5) one concludes that in this case we expect heavy gluinos, typically in the TeV range. As for the number of surviving $\chi'$s in this case, what is crucial is whether $m_\chi$ is larger or smaller than $M_W$. Indeed, for $m_\chi > M_W$ the annihilation channels $\chi\chi \rightarrow WW, ZZ, t\bar{t}$ reduce $\Omega_\chi$ too much. If $m_\chi < M_W$ then acceptable contributions of $\chi$ to $\Omega_{CDM}$ are obtainable in rather wide areas of the $(\mu - M_Z)$ parameter space. Once again I emphasize that the case $\chi$ being a pure higgsino is of particular relevance for LEP 200 given that in this case $\chi$ masses up to $M_W$ can be explored.

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 $\chi$ abundance provides a correct $\Omega_{CDM}$ 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 $\chi$ masses up to 150 GeV it is possible to find sizable regions in the SUSY parameter space where $\Omega_\chi$ acquires interesting values for the DM problem. A detailed and updated analysis is presented in ref. [15] where one can compare the allowed SUSY parameters area with or without the constraint $0.1 < \Omega_\chi h^2 < 0.7$, where $h$ is the Hubble parameter.

There is a further phenomenological constraint which helps in restricting even more severely the available regions of SUSY parameter space where $\Omega_\chi h^2$ can be relevant for the DM problem: it is the recent measurement of the decay $b \rightarrow s + \gamma$ at the inclusive level by the CLEO collaboration. Two papers [16] have recently thoroughly investigated the problem of the direct detection of relic neutralinos in processes of neutralino–nucleus scattering including the constraint arising from the experimental result of $BR(b \rightarrow s + \gamma)$. It turns out that large portions of the SUSY parameter space where it would be possible to have a neutralino–nucleus scattering rate high enough to be detectable in the next round of experiments predict very large values for $BR(b \rightarrow s + \gamma)$ vastly exceeding the experimental result. However, there still survive particular regions where rates as high as $10^{-11}$ events/kg/day for a $^{76}Ge$ detector are allowed. This is the case, for instance, for relatively large $\tan\beta$ ($\tan\beta \sim 20$) and moderate values of the SUSY parameters ($\tilde{m} = 200$ GeV, $\mu = -300$ GeV, $M_Z = 100$ GeV). For a complete discussions I
refer the interested reader to the works of ref. [16].

I close this section with a remark concerning the possibility that gauginos are massless, i.e. $M_1 = M_2 = M_3 = 0$, to start with and that $R$–invariance (the continuous $U(1)$ symmetry associated with the fermionic partners of the gauge bosons, not to be confused with the discrete $R$–parity) is broken spontaneously by Higgs $VEV’s$ or else explicitly by dimension 2 or 3 SUSY–breaking terms in the low energy effective lagrangian. Gluino and lightest neutralino masses then depend on only a few parameters. For a breaking scale of a few hundred GeV or less, the gluino and the lightest neutralino have masses typically in the range $10^{-1} – 2 \text{ GeV}$. On the other hand, for a SUSY–breaking scale several TeV or larger, radiative contributions can yield gluino and lightest neutralino masses of $O(50–300) \text{ GeV}$ and $O(10–30) \text{ GeV}$, respectively. As long as the Higgs $VEV’s$ are the only source of $R$–invariance breaking, or if SUSY breaking only appears in dimension 2 terms in the effective lagrangian, the gluino is generically the lightest SUSY particle, hence modifying the usual phenomenology (and in particular the conventional view of the DM in SUSY) in interesting ways. For reasons of space I cannot deal more with this interesting (or at least curious) issue here and I recommend in particular sect. 5 of our paper [17] with G. Farrar for hints at how the DM problem may be affected by the initial presence of a continuous $U(1)$ $R$–symmetry in supergravity models.

6. LEPTON NUMBER VIOLATION IN SUSY

In the previous section I discussed the more conventional SUSY schemes where $R$ parity is imposed to avoid all the $B$ and $L$ violating operators in the superpotential. From the cosmological point of view the most important consequence of the presence of $R$ is that there exists a stable SUSY particle which has good chances to constitute the CDM in an MDM scenario. As for the hot part of the MDM one can think of neutrinos getting a small mass (in the eV range). In some SUSY GUT’s like SO(10) this is naturally achieved through a see–saw mechanism.

Let me comment now the alternative possibility that $R$–parity is replaced by some other symmetry, for instance $B$–parity, allowing for $B$ or $L$ explicit violation in the superpotential. The removal of $R$–parity has an unpleasant consequence for the DM problem: we lose our beloved CDM candidate represented by the stable LSP. In models with broken $R$–parity the LSP can decay into ordinary particles and, generally, these decays are much faster than what would be required to make
the LSP survive until today.

The only exceptions are situations of extremely tiny violations of $R$-parity. An example is offered in ref. [18]. Not only can the lightest neutralino still be the CDM today, but its slow decays can have an experimental impact: for instance, we considered the possibility of the LSP radiative decays into a $\nu + \gamma$ with a possibly “visible” neutrino line. The negative result of a search performed at Kamiokande of such neutrinos led to a sharp improvement [19] on the bounds of the LSP lifetime (it turns out that $\tau_{LSP}$ must exceed the Universe lifetime by several orders of magnitude).

Although the absence of $R$ parity carries the bad news that in general we lose the obvious SUSY candidate for CDM, it can have a positive impact on the other side of a mixed dark matter (MDM) scenario, i.e. it can yield a good amount of HDM. The point is that $R$ violation is accompanied by $L$ violation (for instance in schemes with $B$-parity), hence allowing for nonvanishing neutrino (Majorana) masses. In addition to the presence of $m_\nu$ there are several other important astrophysical implications: possibly large neutrino magnetic moments, new features in the implementation of the MSW mechanism for the solar neutrino problem, etc. [20].

The explicit violation of $L$ through the presence of $L$ violating operators in the superpotential is severely limited not so much by phenomenological constraints [21], but rather by a powerful cosmological argument related to the survival of the cosmic matter–antimatter asymmetry [22].

The argument goes as follows. It is well–known that owing to the anomalous character of the $L$ and $B$ currents, these two numbers are violated at the quantum level. Only the combination $B - L$ is conserved. Although these violations are unlikely to produce any visible effect at zero temperature, they become quite relevant at high temperature [23]: the associated $B$ and $L$ violating processes have rates larger than the expansion rate of the Universe (at least for $100 \text{ GeV} < T < \text{critical temperature of the electroweak phase transition}$, but, presumably, also for $T > T_c$), hence leading to an equal erasure of the pre–existing $B$ and $L$ asymmetries. Hence, if one starts with $\Delta B = \Delta L$, which is the case in GUT’s with $B - L$ conservation, one ends up with $\Delta B = \Delta L = 0$ at the electroweak phase transition.

Whether these same quantum effects which are responsible for the cosmic $\Delta B$ erasure can be used to produce a new $\Delta B$ at the time of the electroweak phase transition is very doubtful. The survival of a lately produced $\Delta B$ seems
to require an excessively light Higgs boson in the SM and also the amount of CP violation is unlikely to be sufficient to obtain a sizeable $\Delta B$. However, both these objections are far from being settled and further work is needed to make some final assessment on this intriguing issue. A safer way to solve problem is represented by a different boundary condition at the GUT scale with $\Delta B \neq \Delta L$. If this is the case, given that quantum effects preserve $B - L$ it is never possible to reach a total erasement of $\Delta B$. This is the reason why models like SO(10) where $B - L$ is violated (hence allowing for $\Delta B \neq \Delta L$) are certainly favoured with respect to GUT’s with $B - L$ conservation (like SU(5)). Moreover SO(10) schemes can lead to neutrino masses in the convenient range to provide viable candidates for HDM.

All what I said above holds provided that during the interval time from the production of the cosmic $\Delta B$ (for example at the GUT time) down to the electroweak phase transition no other $B$ or $L$ violating interaction is in equilibrium apart from the abovementioned anomalous quantum effects. For instance, if $R$ violating processes are present and are fast enough to be in equilibrium at some moment, since they violate either $B$ or $L$ they certainly violate $B - L$ and hence no combination of $\Delta B$ and $\Delta L$ can survive (independently from whether $\Delta B = \Delta L$ or $\Delta B \neq \Delta L$ to start with). Requiring the $R$–violating induced processes to be out of equilibrium places such a severe bound [22] on the strength of the $R$ violation in the superpotential that certainly one could forget about any phenomenological implication of $R$ breaking. As usual, however, this is not the end of the story concerning SUSY models without $R$ parity. Several solutions have been pointed out to let $\Delta B$ survive even in the presence of non–negligible $R$–breaking effects. Nervertheless the above cosmological observation represents a severe warning for the construction of consistent SUSY schemes which are alternative to those with the traditional matter $R$–parity.

One final comment on $R$–parity breaking is in order. We know that many continuous or global symmetries of the initial lagrangian can be spontaneously broken. One might wonder whether $R$–parity can undergo a similar destiny. Long ago it was pointed out [25] that there are regions of the SUSY parameter space where the minimization of the scalar potential leads to a nonvanishing $VEV$ for the scalar partner of the neutrino, the sneutrino. This would correspond to the spontaneous breaking of $L$ and $R$–parity. By now we know that such a breaking is phenomenologically forbidden. Indeed, the $Z$ boson could decay into the Goldstone boson associated to the breaking of $L$ and the scalar partner of it. The stringent
bound on the invisible width of the $Z$ excludes this possibility.*

Alternatively one can supplement the usual particle spectrum of the minimal SUSY model with one or more gauge singlet scalar superfields which carry $L$ and acquire a $V_{EV}$ \[^{28,29}\]. In this case the Goldstone boson being a gauge singlet does not couple to the $Z$ boson. In relation to the above considerations on baryogenesis and $R$-breaking, it is relevant to notice that the breaking of $R$ can be induced radiatively, i.e. by the evolution of the singlet masses dictated by the renormalization group equations. It was recently shown \[^{29}\] that this radiative breaking can delay the breaking of $R$ down to temperature so low that the $B$ violating quantum effects are no longer effective, i.e. typically $T < 100 \text{ GeV}$.

### 7. MIXED AND WARM DM

As discussed in the Introduction, schemes with pure hot DM or pure cold DM seem disfavoured by recent (and also less recent) observations. Among the new options which are presently envisaged I think that the following two are of particular interest for particle physicists: mixed DM (MDM) and warm DM.

MDM \[^{30}\] relies on a scenario where $\Omega_{CDM} \simeq 2\Omega_{HDM} \simeq 0.6$ and $\Omega_B \lesssim 0.1$. In principle one does not have to sweat so much to realize a scheme of this kind. Take a SUSY model with $R$—parity where neutrinos are massive. Then the lightest neutralino can play the role of CDM, while a neutrino of few $eV$s yields the HDM. Choosing the parameters conveniently one can obtain the prescribed cocktail of C- and H- DM. The problem that I see is just in this convenient choice of parameters. This is another way to say that one actually performs a fine–tuning to obtain the correct amount of $\Omega_{CDM}$ and $\Omega_{HDM}$ and this is certainly unsatisfactory. This is the reason which prompted some authors to investigate some possible common origin for HDM and CDM in order to justify close relation of their contributions to $\Omega$. In the work of ref. \[^{31}\] it was proposed to have the relative abundances of the HDM and CDM components set by the same scale. In their model, this is

* It was recently discussed the possibility that gravitational effects spoil any global symmetry \[^{26}\]. If this is the case, $L$ might be explicitly broken very tinily. The subsequent “spontaneous” breaking through a $V_{EV}$ of the sneutrino gives rise to a pseudo–Goldstone boson. Interestingly enough, even though the explicit breaking is very small, the mass of this particle can easily exceed the $Z$ mass hence preventing the abovementioned decay which contributed to the $Z$ invisible width \[^{27}\].
the scale of B-L spontaneous breaking of $O(1 \text{ TeV})$. The HDM is given by the tau neutrino, while CDM is provided by the fermionic partner of the Goldstone boson associated to the B-L breaking.

Together with Bonometto and Gabbiani, we proposed \cite{32} an example where one same particle may play the twofold role of HDM and CDM. In SUSY the axion possesses a fermionic partner, the axino ($\tilde{a}$). In fact, the $\tilde{a}$ is likely to be the lightest SUSY particle. Now, axinos can be produced via two entirely different different mechanisms in these models. First there are the axinos which are produced with the axions and were formerly in thermal equilibrium with the other components of the Universe, subsequently decoupling at a temperature $< V_{PQ}$ (the Peccei–Quinn scale) much higher than their mass. This $\tilde{a}$ component will be an effective CDM as only fluctuations involving masses $\lesssim 0.1 M_\odot$ will be erased at its derelativization. It was shown that they can account for $\Omega$ close to one \cite{33}. This kind of “primordial” axinos are not the only axinos surviving today. Indeed if the $\tilde{a}$ is the lightest SUSY particle, all the SUSY particle must eventually decay into it.

Calling $\chi$ the lightest neutralino, we can expect the typical decay $\chi \to \tilde{a} + \gamma$ to occur through a supersymmetrization of the ordinary $a - \gamma \gamma$ coupling.

These “second hand” axinos can easily behave as hot dark matter, derelativizing at a redshift $z \sim 10^4$. Accordingly, fluctuations in such component will be erased up to a mass $\sim 10^{15} M_\odot$.

The detailed study of the conditions which make this scheme a viable MDM scenario is presented in ref. \cite{32}. The major ingredients are a Peccei–Quinn scale of $O(10^{10} \text{ GeV})$, heavy sfermions in the TeV range and the lightest neutralino being a pure gaugino.

An interesting alternative to MDM is the presence of just one DM particle which is neither cold nor hot. This warm candidate may be represented for instance by a sterile neutrino which is somewhat heavier but less abundant than the usual HDM neutrinos. Clearly one must be very careful about the contribution of these extra degrees of freedom at the time of nucleosynthesis (they must contribute less than the equivalent of half a neutrino species). The essential point of warm DM is that it can reduce the damping scale corresponding to the free–streaming distance that was previously introduced. If for an ordinary HDM neutrino this damping scale is of $O(10^{15} M_\odot)$, for the kind of warm sterile neutrinos discussed in ref. \cite{34} this is lowered to $10^{13} M_\odot$ hence increasing the power on smaller scales (typically scales 1-5 Mpc).
Another example of warm DM candidate results from the “spontaneous” breaking of a quasi-exact $L$ symmetry (as explained in the previous footnote). A pseudo-Goldstone boson with a mass in the keV range and with tiny interaction with ordinary matter has been shown [35] to be a suitable candidate for warm DM.

All these attempts of a mixed and warm DM to realize a better fit to data at different scales are certainly interesting. However I must confess that my overall impression is that we are far from having an appealing scenario with some compelling reason from the particle physics point of view. In this respect scenarios with pure CDM or pure HDM were much more attractive. The “canonical” final sentence that more work is needed definitely applies very well to the present situation in this field.

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