THROUGH THE LOOKING-GLASS:
ALICE'S ADVENTURES IN MIRROR WORLD [*]

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We briefly review the concept of a parallel 'mirror' world which has the same particle physics as the observable world and couples to the latter by gravity and perhaps other very weak forces. The nucleosynthesis bounds demand that the mirror world should have a smaller temperature than the ordinary one. For this reason its evolution should substantially deviate from the standard cosmology as far as the crucial epochs like baryogenesis, nucleosynthesis etc. are concerned. In particular, we show that in the context of certain baryogenesis scenarios, the baryon asymmetry in the mirror world should be larger than in the observable one. Moreover, we show that mirror baryons could naturally constitute the dominant dark matter component of the Universe, and discuss its cosmological implications.

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[*] For technical reasons we did not manage to put this contribution in Chapter 3: Unification, Extra Dimensions, Cosmology, to which it belongs — Editors.

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Ian's death was one of the most shocking events which I have ever experienced. I met him for the first time 25 years ago; we were young and just making our first steps in physics. Karen Avetovich Ter-Martirosyan, gabbling in his typical manner, presented to me a smiling handsome boy with long black hair and beautiful long eyelashes: "This is Ian Kogan, very promising, very strong, brilliant ..." – obviously KA admired him. Since then we were close friends. Several nights we spent together with a good bottle of wine discussing anything in physics and beyond physics, but only in recent years did we start to collaborate on concrete projects – in particular, we both were intrigued by the idea of parallel worlds: I was motivated by the problems of dark matter and baryogenesis, and he was attracted by bi-gravity. Just a day before his death he sent an e-mail message from Trieste, amusing as always – "So what old pirate, did you start to write the draft? Italy is wonderful — Ahhh! I leave tomorrow and then I'll call you up from Oxford ... " He was going to visit me in Gran Sasso in a couple of weeks.

Ian's imagination and attraction to "fantastic" ideas were beautifully complemented with a wonderful sense of humor. Definitely there was something of Lewis Carroll's in him ... One of his papers is devoted to the Alice strings – possible passages to the parallel worlds ...

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'Now, if you'll only attend, Kitty, and not talk so much, I'll tell you all
my ideas about Looking-glass House. First, there's the room you can see
through the glass – that's just the same as our drawing-room, only the things
go the other way... the books are something like our books, only the words
go the wrong way: I know that, because I've held up one of our books to
the glass, and then they hold up one in the other room. I can see all of it
when when I get upon a chair – all but the bit just behind the fireplace. I
do so wish I could see that bit! I want so to know whether they've a fire in
the winter: you never can tell, you know, unless our fire smokes, and then
smoke comes up in that room too – but that may be only pretence, just to
make it look as if they had a fire...

'How would you like to leave in the Looking-glass House, Kitty? I wander
if they'd give you milk in there? Perhaps Looking-glass milk isn't good to
drink – but Oh, Kitty! Now we come to the passage. You can just see a
little peep of the passage in Looking-glass House, if you leave the door of
our drawing-room wide open: and it's very like our passage as far as you
can see, only you know it may be quite on beyond. Oh, Kitty, how how nice
it would be if we could get through into Looking-glass House! Let's pretend
there's a way of getting through into it, somehow, Kitty... Why, it's turning
into a sort of mist now, I declare! It'll be easy enough to get through – 'She
said this, though she hardly knew how she had got there...

In another moment Alice was through the glass, and had jumped lightly
down into the Looking-glass room. The very first thing Alice did was to look
whether there was a fire in the fireplace, and she was quite pleased to find
that there was a real one, blazing away as brightly as the one she had left
behind. 'So I shall be as worm here as I was in the room,' thought Alice:
'warmer, in fact, because there'll be no one here to scold me away from the
fire. Oh, what fun it'll be, when they see me through the glass in here, and
ca'n't get at me!'

Lewis Carroll, *"Through the Looking-Glass"*

1. Introduction

Lewis Carroll's *Alice* probably was the first who seriously considered that
the world beyond the mirror – *"Looking-Glass House"* – is real: "just same
as our world, only the things go other way..." Observable elementary par-
ticles have left-handed $(V - A)$ weak interactions which violate P-parity in
the strongest possible way. However, there could exist a hidden mirror world

of particles, an exact copy of our world, only that mirror particles experience right-handed $(V + A)$ weak interactions. Each ordinary particle has its mirror twin: "the mirror particles are something like our particles, only the chiralities go the wrong way..." Such a duplication of the worlds would restore the left-right symmetry of Nature, as was suggested by Lee and Yang [2]. The phenomenological implications of such a parallel world were first addressed by Kobzarev, Okun and Pomeranchuk, who also introduced the term "Mirror World" [3], and several other papers followed [4]- [12].

The basic concept is to have a theory given by the product $G \times G'$ of two identical gauge groups with identical particle contents, which could naturally emerge e.g. in the context of the $E_8 \times E_8'$ superstring. Once the gauge factor G describes interactions of observable particles: quarks and leptons, Higgses, etc., then its gauge counterpart G' describes the world with analogous particle content: mirror quarks and leptons, mirror Higgses, etc. (From now on all fields and quantities of the mirror (M) sector will be marked by $'$ to distinguish them from the ones belonging to the observable or ordinary (O) world.) The M-particles are singlets of G and vice versa, the O-particles are singlets of G'. A discrete symmetry $G \leftrightarrow G'$ interchanging corresponding fields of G and G', mirror parity, guarantees that the two particle sectors have identical Lagrangians, with all coupling constants (gauge, Yukawa, Higgs) having the same pattern, and thus their microphysics is identical.[a] The two worlds communicate through gravity, but there are also other possible ways.[b]

If the mirror sector exists, then the Universe along with the ordinary photons, neutrinos, baryons, etc. should contain their mirror partners. One could naively think that due to mirror parity the O- and M- particles should have the same cosmological abundances and hence the two sectors should have the same cosmological evolution. However, this would be in the immediate conflict with the Big Bang nucleosynthesis (BBN) bounds on the effective number of extra light neutrinos, since the mirror photons, electrons and neutrinos would give a contribution to the Hubble expansion rate equivalent to $\Delta N_\nu \simeq 6.14$. Therefore, in the early Universe the M-system should have a lower temperature than ordinary particles. This situation is plausible if the following conditions are fulfilled:

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[a] Parity between two worlds can be spontaneously broken, e.g. the electroweak symmetry breaking scales in two sectors can be different, which would lead to somewhat different particle physics in mirror sector [13, 14, 16].

[b] For example, ordinary photons could have kinetic mixing with mirror photons [7–9], ordinary neutrinos could mix with mirror neutrinos [12, 13], two sectors could have a common gauge symmetry of flavor [15] or common Peccei–Quinn symmetry [16].

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A. After the Big Bang (post-inflationary reheating) the two systems get different initial temperatures, namely the temperature in the M-sector is lower than in the visible one, $T' < T$. This can be naturally achieved in certain models of inflation [14, 17, 18].

B. The two systems interact very weakly, so that they do not come into thermal equilibrium with each other during the Universe expansion. This condition is automatically fulfilled if the two worlds communicate only via gravity. If there are some other effective couplings between O- and M- particles, they have to be properly suppressed.

C. Both systems expand adiabatically, without significant entropy production.

If these conditions are satisfied, two sectors with different initial temperatures, evolving independently during the cosmological expansion, maintain the ratio of their temperatures T'/T nearly constant at later stages. In this way, if $T'/T \ll 1$, mirror sector would not affect primordial nucleosynthesis in the ordinary world.

At present, the temperature of ordinary relic photons is $T \approx 2.75$ K, and the mass density of ordinary baryons constitutes about 5% of the critical density. Mirror photons should have smaller temperature $T' < T$, so their number density is $n'_\gamma = x^3 n_\gamma$, where $x = T'/T$. This ratio is a key parameter in our further considerations insofar as it remains nearly invariant during the expansion of the Universe. The BBN bound on ΔN_ν implies the upper bound $x < 0.64 \, \Delta N_\nu^{1/4}$. As for mirror baryons, *ad hoc* their number density n'_b can be larger than n_b, and if the ratio $\beta = n'_b/n_b$ is about 5 or so, they could constitute the dark matter of the Universe.

In this paper we discuss the cosmological implications of the mirror sector. We show that due to the temperature difference, in the mirror sector all key epochs such as those of baryogenesis, nucleosynthesis, etc. proceed at somewhat different conditions compared to the observable Universe. In particular, we show that in certain baryogenesis scenarios the M-world gets a larger baryon asymmetry than the O-sector, and it is pretty plausible that $\beta > 1$ [19]. This situation emerges in a particularly appealing way in the leptogenesis scenario due to the lepton number leaking from the O- to the M-sector which leads to $n'_b \geq n_b$, and can thus explain the near coincidence of visible and dark components in a rather natural way [20].

In discussing the physics and the cosmology of the two worlds, let us introduce two observers: Ordinary observer *Olga* and Mirror observer *Maxim*.[c]

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[c] I do not specify here why I have chosen these names, but Ian Kogan would know.

The world of *Maxim* is a hidden sector for *Olga* and vice versa, the world of *Olga* is a hidden world for *Maxim*. Could they by some experimental and theoretical means deduce the existence of the parallel hidden sectors? Also *Alice* can be involved as a super-observer who can see the whole theory, and as a possible messenger between two worlds.

In the presently popular language of extra dimensions and brane-worlds, the concept of two parallel worlds can be visualized in a simple way. One could consider e.g. a five-dimensional theory with compactified and orbifolded fifth dimension (S_1/Z_2) with parallel 3D-branes located in two fixed points, so that the ordinary matter is localized on the left-brane L and the mirror matter is localized on the right-brane R. In this view, one would simply say that *Olga* lives on the L-brane and *Maxim* on the R-brane, while *Alice* can propagate in the bulk.

2. Mirror world and mirror symmetry

Let us now discuss in more detail *Alice*'s theory of two parallel worlds. We consider two identical gauge factors $G \times G'$ with identical representations. Mirror parity is understood as a discrete symmetry under $G \to G'$, when all ordinary particles (fermions, Higgses and Gauge fields) exchange places with their mirror partners ('primed' fermions, Higgses and Gauge fields), so that the Lagrangian of the O-sector

$$\mathcal{L} = \mathcal{L}_{\text{Gauge}} + \mathcal{L}_{\text{Higgs}} + \mathcal{L}_{\text{Yuk}} \tag{1}$$

transforms into the Lagrangian of M-sector,

$$\mathcal{L}' = \mathcal{L}'_{\text{Gauge}} + \mathcal{L}'_{\text{Higgs}} + \mathcal{L}'_{\text{Yuk}} \tag{2}$$

and the whole Lagrangian $\mathcal{L}' + \mathcal{L}$ remains invariant.

Let us consider, for simplicity, that the O-world is described by the Standard Model (SM) based on the gauge symmetry $G_{\text{SM}} = SU(3) \times SU(2) \times U(1)$, which has a chiral fermion content with respect to the electroweak gauge factor $SU(2) \times U(1)$. Fermions are represented as Weyl spinors, the left-handed (L) quarks and leptons f_L transforming as doublets and the right-handed (R) ones f_R as singlets,[d]

$$f_L: \quad q_L = \begin{pmatrix} u_L \\ d_L \end{pmatrix}, \quad l_L = \begin{pmatrix} \nu_L \\ e_L \end{pmatrix}; \qquad f_R: \quad u_R, \ d_R, \ e_R. \tag{3}$$

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[d] Here and in the following, we omit the family indices for simplicity.

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Then the field operators $\tilde{f}_R = C\gamma_0 f_L^*$ and $\tilde{f}_L = C\gamma_0 f_R^*$, C being the charge conjugation matrix, describe antifermions which have opposite gauge charges as well as opposite chiralities with respect to what we call fermions,

$$\tilde{f}_R: \quad \tilde{q}_R = \begin{pmatrix} \tilde{u}_R \\ \tilde{d}_R \end{pmatrix}, \quad \tilde{l}_R = \begin{pmatrix} \tilde{\nu}_R \\ \tilde{e}_R \end{pmatrix}; \qquad \tilde{f}_L: \quad \tilde{u}_L, \; \tilde{d}_L, \; \tilde{e}_L \,. \tag{4}$$

In addition, we prescribe a global baryon charge $B = 1/3$ to quarks q_L, u_R, d_R, (so that baryons consisting of three quarks have $B = 1$), and a lepton charge $L = 1$ to the leptons l_L, e_R. Hence antiquarks $\tilde{q}_R, \tilde{u}_L, \tilde{d}_l$ have $B = -1/3$, and antileptons $\tilde{l}_R, \tilde{e}_L$ have $L = -1$.

The Gauge and Higgs parts in the Lagrangian $\mathcal{L}$ are self-explanatory, while the fermion Yukawa couplings with the Higgs doublet(s) $\phi = \phi_{u,d}$ can be conveniently presented in the form [e]

$$\mathcal{L}_{\text{Yuk}} = \mathcal{W} + \mathcal{W}^\dagger \,, \tag{5}$$

$$\mathcal{W} = \tilde{f}_L Y f_L \phi \equiv \tilde{u}_L Y_u q_L \phi_u + \tilde{d}_L Y_d q_L \phi_d + \tilde{e}_L Y_e l_L \phi_d \,,$$

$$\mathcal{W}^\dagger = f_R Y^* \tilde{f}_R \tilde{\phi} \equiv u_R Y_u^* \tilde{q}_R \tilde{\phi}_u + d_R Y_d^* \tilde{q}_R \tilde{\phi}_d + e_R Y_e^* \tilde{l}_R \tilde{\phi}_d \,,$$

where $Y = Y_{u,d,e}$ are the Yukawa constant matrices and $\tilde{\phi}_{u,d} \equiv \phi_{u,d}^*$ (C-matrix and the sign of transposition are omitted for simplicity). Therefore, $\mathcal{W}$ is a holomorphic function of the L-fields and $\mathcal{W}^*$ of R-fields.

On the other hand, the physics of the M-sector, based on the mirror Standard Model with the gauge symmetry $G'_{\text{SM}} = SU(3)' \times SU(2)' \times U(1)'$ has the analogous content of fermion fields:

$$f'_L: \quad q'_L = \begin{pmatrix} u'_L \\ d'_L \end{pmatrix}, \quad l'_L = \begin{pmatrix} \nu'_L \\ e'_L \end{pmatrix}; \qquad f'_R: \quad u'_R, \; d'_R, \; e'_R \,, \tag{6}$$

and that of anti-fermions

$$\tilde{f}'_R: \quad \tilde{q}'_R = \begin{pmatrix} \tilde{u}'_R \\ \tilde{d}'_R \end{pmatrix}, \quad \tilde{l}'_R = \begin{pmatrix} \tilde{\nu}'_R \\ \tilde{e}'_R \end{pmatrix}; \qquad \tilde{f}_L: \quad \tilde{u}'_L, \; \tilde{d}'_L, \; \tilde{e}'_L \,. \tag{7}$$

For definiteness, let us prescribe mirror fermion numbers: $B' = 1/3$ to quarks q'_L, u'_R, d'_R and $L' = 1$ and leptons l'_L, e'_R, and so antiquarks $\tilde{q}'_R, \tilde{u}'_L, \tilde{d}'_l$ have $B' = -1/3$, and antileptons $\tilde{l}'_R, \tilde{e}_L$ have $L' = -1$.

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[e] In the minimal SM, ϕ_u and ϕ_d simply are conjugated fields, $\phi_d \sim \phi_u^*$. However, we keep in mind that generally in the extensions of the SM, in particular, in the supersymmetric extension, ϕ_u and ϕ_d are independent ("up" and "down") Higgs doublets with different vacuum expectation values (VEV) $\langle\phi_u\rangle = v_u$ and $\langle\phi_d\rangle = v_d$, and their ratio is known as $\tan\beta = v_u/v_d$.

Yukawa couplings have the form analogous to (5),

$$\mathcal{L}'_{\text{Yuk}} = \mathcal{W}' + \mathcal{W}'^{\dagger}, \tag{8}$$
$$\mathcal{W}' = \tilde{f}'_L Y' f'_L \phi' \equiv \tilde{u}'_L Y'_u q'_L \phi'_u + \tilde{d}'_L Y'_d q'_L \phi'_d + \tilde{e}'_L Y'_e l'_L \phi'_d,$$
$$\mathcal{W}'^{\dagger} = f'_R Y'^* \tilde{f}'_R \tilde{\phi}' \equiv u'_R Y'^*_u \tilde{q}'_R \tilde{\phi}'_u + d'_R Y'^*_d \tilde{q}'_R \tilde{\phi}'_d + e'_R Y'^*_e \tilde{l}'_R \tilde{\phi}'_d,$$

where $\phi' = \phi'_{u,d}$ are the mirror Higgs doublets.

What kind of discrete symmetries can such a theory have?

Let us first consider the O-sector separately. The weak interactions of ordinary particles break one of the possible fundamental symmetries of the Nature, parity, in the strongest possible way. The physics is not invariant under the coordinate transformation $x \to -x$, and hence left- and right-handed systems of the coordinates are not equivalent.

Namely, for what she calls particles, baryons and leptons, (probably because she herself is made up of them), the weak interactions have the left-handed $(V - A)$ form.

The particle content of the Standard Model and hence its Lagrangian is not symmetric under the exchange of the L and R particles, $f_L \leftrightarrow f_R$. In particular, the gauge bosons of $SU(2)$ couple to the f_L fields but do not couple to f_R ones. In fact, in the limit of unbroken $SU(2) \times U(1)$ symmetry, f_L and f_R are essentially independent species with different quantum charges. The only reason why we call e.g. two Weyl fermions $e_L \subset l_L$ and e_R respectively as the left- and right-handed electrons is that after the electroweak breaking down to $U(1)_{\text{em}}$ these two have the same electric charges and form a massive Dirac fermion $\psi_e = e_L + e_R$.

There exist left-right extensions of the Standard Model, with the electroweak gauge symmetry extended to $SU(2)_L \times SU(2)_R \times U(1)_{B-L}$. Now f_L are doublets of $SU(2)_L$ and f_R (including right-handed neutrinos) are doublets of $SU(2)_R$. The Lagrangian of such a theory can be invariant under the exchange $f_L \to f_R$ if at the same time the two gauge sectors interchange places, $SU(2)_L \to SU(2)_R$. However, experiment tells us that if such a parity exists, than it should be spontaneously broken; the mass of the W_R gauge bosons should be much larger than the mass of ordinary W_L bosons.

Once again, for what we call particles f (3), their weak interactions are left-handed $(V - A)$, while in terms of antiparticles $\tilde{f}$ (4), the weak interactions would be seen as right-handed $(V + A)$, since now only R states couple to the $SU(2)$ bosons. In the context of the SM or its extensions, the symmetry between particles f and antiparticles $\tilde{f}$, CP-parity, could be the the only exact fundamental symmetry between the left and right: $f_L \to \tilde{f}_R$,

$(f_R \to \tilde{f}_L)$, as far as it is respected by the gauge interactions. But not by the Yukawa terms! Although the terms (5) are presented in an rather symmetric manner between L and $\tilde{R}$ fields, $\mathcal{W}$ being a holomorphic function of only L fields and $\mathcal{W}^*$ of $\tilde{R}$ fields, they are not invariant under the interchange $L \to \tilde{R}$ due to irremovable complex phases in the Yukawa coupling matrices $Y_{u,d,e}$. As we know, after the electroweak symmetry breaking, the complexity propagates to the CKM matrix of the quark mixing and leads to observable CP-violation effects. Hence, neither CP-parity is respected. Nature once again demonstrates that "the only good parity is a broken parity".

Clearly, one could always redefine the notion of particles and antiparticles, to rename particles as antiparticles and vice versa. The natural choice for what to call particles is given by the content of matter in our Universe. Matter, at least in our galaxy and its wide neighborhood, consists of baryons q while antibaryons $\tilde{q}$ can be met only in accelerators or perhaps in cosmic rays. However, if by chance we lived in the antibaryonic island of the Universe, we would claim that our weak interactions are right-handed. Therefore, not only our microphysics but also our Universe does not respect the exchange between matter and antimatter.

The CP-violating effects in particle physics are tiny. In fact, most of the elementary processes are exactly the same between the particles and antiparticles, and CP-violation is observed only in some rare processes, as e.g. leptonic asymmetry in K_L decay.

On the other had, our Universe, presently being completely dominated by matter over antimatter, $n_b/n_{\bar{b}} \gg 10^{10}$, at the early stages was almost symmetric, with a tiny excess of baryons over antibaryons $n_b/n_{\bar{b}} = 1 + \mathcal{O}(10^{-10})$. It is a very profound question, who has prepared our Universe (and how) in the initial state to provide a tiny excess of baryons over antibaryons, and therefore fixed a priority of the $V - A$ form of the weak interactions over the $V + A$ one? It is appealing to think that the baryon asymmetry itself emerges due to the tiny CP-violating features in particle interactions, and it is related to some fundamental physics beyond the Standard Model which is responsible for primordial baryogenesis. Without the CP-violating effects, probably, no Baryon Asymmetry would be possible at all, and the Universe would consist only of light.

Concluding, the particle physics of the ordinary world is not symmetric between the left and right – neither $P : f_L \to f_R$ nor $CP : f_L \to \tilde{f}_R$ are exact symmetries – they are explicitly (or perhaps spontaneously) broken.

However, the whole theory describing both worlds, can be symmetric between left and right. Indeed, consider the transformation of coordinates

P: $x \to -x$ accompanied by the transformation which interchanges O-fields f, ϕ with the corresponding M-fields $\tilde{f}, \tilde{\phi}$ in the following manner:

$$f_L, \tilde{f}_L \to \gamma_0 \tilde{f}'_R, \gamma_0 f'_R, \quad \tilde{f}_R, f_R \to \gamma_0 f'_L, \gamma_0 \tilde{f}'_L, \quad \phi \to \tilde{\phi}' \tag{9}$$

while the gauge fields of G properly transform into their partners in G'. Such a symmetry can be called matter parity MP. Obviously, it implies that the gauge couplings are exactly the same in the two sectors, the Higgs potentials are identical, while for the Yukawa coupling constants we have

$$Y'_{u,d,e} = Y^*_{u,d,e} \,. \tag{10}$$

In this way, the introduction of the M-world does not introduce any new parameter. So, in this case the particle physics of the M-world will be exactly the same in terms of the R-fields $\tilde{f}'_R, f'_R$ as that of the O-world in terms of L-fields $f_L, \tilde{f}_L$. Hence, MP restores the left-right symmetry as a symmetry between the two sectors.[f]

The generalization for the extensions of the Standard Model related to supersymmetry and grand unification is straightforward. Consider a generic supersymmetric gauge theory with $G \times G'$ symmetry where the O-sector is presented by a set of left-chiral superfields L in certain representations of G and the M-sector by a set of left-chiral superfields $\tilde{L}$ in analogous (anti)representations of G', so that L are singlets of G' and vice versa, $\tilde{L}$ are singlets of G. One can also explicitly write their conjugated right-chiral superfields $\tilde{R} = L^*$, $\tilde{R}' = L'^*$

For example, in the context of the MSSM, both O-matter and M-matter are presented in terms of left-chiral superfields L and their conjugated right-chiral superfields $\tilde{R}$,[g]

$$\begin{aligned}
L : & \quad q, l, \tilde{u}, \tilde{d}, \tilde{e}, \ \phi_{u,d}; \quad & \tilde{R} : & \quad \tilde{q}, \tilde{l}, u, d, e, \ \tilde{\phi}_{u,d}, \\
L' : & \quad q', l', \tilde{u}', \tilde{d}', \tilde{e}', \ \phi'_{u,d}; \quad & \tilde{R}' : & \quad \tilde{q}', \tilde{l}', u', d', e', \ \tilde{\phi}'_{u,d}
\end{aligned} \tag{11}$$

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[f] Alternatively, one could impose the D-symmetry under transformation $f_{L,R} \leftrightarrow f'_{L,R}$, $\phi \leftrightarrow \phi'$, which would imply $Y'_{u,d,e} = Y_{u,d,e}$. This case is nothing but direct duplication. (Obviously, if both MP and D are imposed, the whole theory would be CP-invariant.) Either type of parity implies that the two sectors have the same particle physics. If the two sectors are separate and do not interact by forces other than gravity, the difference between D and MP is rather symbolic and does not have any profound implications. However, in scenarios with some particle messengers between the two sectors this difference can be important and can have dynamical consequences.

[g] In the context of $N = 1$ supersymmetry, fermions and Higgses become chiral superfields, and formally they can be distinguished only by matter parity.

with the superpotential terms

$$W = \tilde{u}Y_u q\phi_u + \tilde{d}Y_d q\phi_d + \tilde{e}Y_e l\phi_d \; + \; \mu\phi_u\phi_d \,,$$
$$W' = \tilde{u}'Y_u' q'\phi_u' + \tilde{d}'Y_d' q'\phi_d' + \tilde{e}'Y_e' l'\phi_d' + \mu'\phi_u'\phi_d' \,. \tag{12}$$

In $SU(5) \times SU(5)'$ model L should include the fermion superfields $\bar{5} + 10$ and the Higgs superfields $5, \bar{5}$ and 24 of $SU(5)$, and $\tilde{L}$ the same superfields of $SU(5)'$.

Then the Lagrangian reads

$$\mathcal{L}_{\mathrm{mat}} = \int d^2\theta W(L) + \int d^2\overline{\theta} W^*(\tilde{R}) + \mathcal{L}_{\mathrm{Gauge}} \,,$$
$$\mathcal{L}_{\mathrm{mat}}' = \int d^2\theta W'(L') + \int d^2\overline{\theta} W'^*(\tilde{R}') + \mathcal{L}_{\mathrm{Gauge}}' \tag{13}$$

with superpotentials $W(L)$ and $W'(L')$ being holomorphic functions respectively of the superfields L and $\tilde{L}$.

Then, M-parity can be understood as the transformation of all left superfields L of G into corresponding right superfields $\tilde{R}'$ of G', $L \to \tilde{R}'$ and $L' \to \tilde{R}$, accompanied by appropriate exchange between the vector superfields of G and G'; hence it implies that W and W'^* have the same functional shape, or W and W' are complex conjugated. Therefore, all complex coupling constants in W' should have the opposite phase with respect the corresponding ones in W. In particular, for the superpotential (13) MP implies that $Y_{u,d,e}' = Y_{u,d,e}^*$ and $\mu' = \mu^*$.

Concluding, the MP implies that the M-world has the same physics in terms of the R-states as the ordinary one in terms of the L-states. However, this does not mean that macroscopic realizations of M-worlds would necessarily be the mirror reflection of O-world. The sign of baryon asymmetry is crucial for determination of the nature of M-world.

Let us discuss now how the particle physics of the two worlds can be seen by ordinary and mirror observers in different situations.

In the ordinary world baryons dominate over antibaryons, i.e. $B > 0$. In this view, the ordinary observer *Olga* identifies f species as particles (matter) and $\tilde{f}$-species as antiparticles (antimatter). Her experimental results can be then formulated as:

- P is broken. Matter of O-world has left-handed nature: weak interactions of O-particles have $V - A$ form, neutrinos are L-handed.
- CP is broken. Decays of K_L meson demonstrate a tiny excess of the positrons e^+ over the electrons e^-.

As for the M-sector, two different situations are possible. The experimental results of the mirror observer *Maxim* would depend on the sign of baryon asymmetry in M-world. A priori it can be either positive or negative. Namely, if in M-world $\tilde{f}'$ species dominate over f', i.e. $B' < 0$, then *Maxim* would identify the former as particles and the latter as antiparticles. In this case his conclusions would be:

- P is broken. Matter of M-world has right-handed nature: weak interactions of O-particles have $V + A$ form, neutrinos are R-handed.
- CP is broken. Decays of the K_L meson demonstrate a tiny excess of positrons e^+ over electrons e^-.

Obviously, this situation will necessarily be realized if baryon asymmetries in the two sectors are induced separately, by the same particle physics mechanism related to CP-violating phases which are opposite for the Yukawa couplings of L and L'.

However, it is also possible that the M-world is dominated by f' species over $\tilde{f}'$, i.e. $B' > 0$. In this case naturally *Maxim* would identify particles as f' and antiparticles as $\tilde{f}'$, and thus he would conclude that:

- P is broken. Matter of O-world has left-handed nature: weak interactions of O-particles have $V - A$ form, neutrinos are L-handed.
- CP is broken. Decays of the K_L meson demonstrate a tiny excess of electrons e^- over positrons e^+.

As we shall see later, such a situation can be realized if baryon asymmetries in both worlds arise by a unique mechanism related to the particle interactions between two worlds via some messengers.

It is not clear how two observers could communicate to each other the information about the microscopic nature of their worlds – can they encode it e.g. into polarized gravitational waves or even photons? (if there is a kinetic mixing between two photons). Perhaps *Alice* could play the role of the messenger: Ian Kogan with collaborators have considered Alice string as a way to travel from one parallel sector to another [22].

3. Neutrinos as messengers between O- and M-worlds

In the context of the minimal Standard Model neutrinos can get Majorana masses from the dimension 5 operator cutoff by a large mass scale M,

$$\frac{1}{2M} (\phi l) A (l \phi) + \text{h.c.} , \tag{14}$$

where A is a symmetric 3×3 matrix of the coupling constants and $\phi \equiv \phi_u$. After substituting the Higgs VEV $\langle \phi \rangle = v$, one obtains the neutrino mass matrix $\hat{m}_\nu = A \cdot v^2/M$. This naturally explains why neutrinos are much lighter than the charged leptons and quarks. The latter are Dirac fermions with masses proportional to the weak scale v, whereas the neutrinos are Majorana fermions with masses $\sim v^2/M$.

A similar operator of dimension 5

$$\frac{1}{2M} (\phi'l')A'(l'\phi') + \text{h.c.} \tag{15}$$

induces the Majorana mass matrix of M-neutrinos, $\hat{m}'_\nu = A' \cdot v'^2/M$, via the M-Higgs VEV $\langle \phi' \rangle = v'$.[h]

However, a mixed operator of dimension 5 is also allowed,

$$\frac{1}{M} \phi l D l' \phi' + \text{h.c.} \,. \tag{16}$$

After substituting the Higgs VEVs it leads to a mixed mass matrix $\hat{m}_{\nu\nu'} = D \cdot vv'/M$, with D being a 3×3 matrix of the coupling constants. Thus, the total 6×6 mass matrix of ordinary neutrinos $\nu \subset l$ and their mirror partners $\nu' \subset l'$ reads [13]

$$\begin{pmatrix} \hat{m}_\nu & \hat{m}_{\nu\nu'} \\ \hat{m}_{\nu\nu'}^T & \hat{m}_{\nu'} \end{pmatrix} = \frac{v^2}{M} \begin{pmatrix} A & \zeta D \\ \zeta D^T & \zeta^2 A' \end{pmatrix} \,, \tag{17}$$

where $\zeta = v'/v$. Mirror Parity (9) imposes the following constraints on the coupling constant matrices

$$A' = A^*, \quad D = D^\dagger \,. \tag{18}$$

In general, this matrix describes six mass eigenstates of Majorana neutrinos, which are superpositions of three ordinary neutrinos and three mirror neutrinos. In the language of neutrino physicists, the O-neutrinos ν_e, ν_μ, ν_τ are *active* neutrinos while the M-neutrinos $\nu'_e, \nu'_\mu, \nu_\tau$ are *sterile* neutrinos. Thus, this model provides a simple explanation of why sterile neutrinos could be light (on the same grounds as the active neutrinos) and could have significant mixing with the active neutrinos.

If parity between the two worlds is exact, $\zeta = 1$, then the neutrinos of the two sectors will be strongly mixed. It seems difficult to reconcile this situation with the present experimental and cosmological limits on the active-sterile neutrino mixing, however it is still premature to conclude that

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[h] As we discuss later, MP can spontaneously broken and in general v' can be different from v.

it is ruled out. If instead mirror parity is spontaneously broken, so that e.g. $\zeta = v'/v \gg 1$, then the active-sterile mixing angles should be small, $\theta_{\nu\nu'} \sim 1/\zeta$, while their mass ratios (per each flavor) scale as $m_\nu/m_{\nu'} \sim 1/\zeta^2 \gg 1$.

The situation when the O- and M-neutrinos have separate mass terms and thus there are no active-sterile neutrino oscillations corresponds to a particular case $D = 0$, and in general it should be motivated by some additional symmetry reasons.

Another interesting case corresponds to $A, A' = 0$ but $D \neq 0$. It can emerge if two sectors have a common lepton number (or $B - L$) symmetry: ordinary leptons l have lepton charges $L = 1$ and mirror ones l' have $L = -1$. Obviously, this symmetry would forbid the terms (14) and (15), while the operator (16) is allowed. Then, 'Majorana' mass terms are absent both for O- and M-neutrinos in (17) and neutrinos are Dirac fermions with *naturally small* masses $\sim vv'/M$, having left components $\nu_L \subset l$ in the O-sector and right components $\tilde\nu'_R \subset \tilde l$ in the M-sector.

Let us consider the situation in the context of the seesaw mechanism; it is nothing but a natural way of generating the effective operators (14) etc. from the renormalizable couplings. In the context of the Standard Model, in addition to the fermions (3) with non-zero gauge charges, one introduces also the gauge singlets, so called right-handed neutrinos $\tilde N_a = \tilde N_{kR}$ (or their conjugated left-handed states $N_a = C\gamma_0\tilde N_a^*$), with the large Majorana mass terms $\frac{1}{2}(M_{ab}N_aN_b + M_{ab}^*\tilde N_a\tilde N_b)$. The mass matrix M is symmetric with respect to indices $a, b = 1, 2, ..., n$, and it is convenient to parametrize it as $M_{ab} = g_{ab}M$, where M is a typical mass scale and g is a matrix of dimensionless Yukawa-like constants.[i] The ordinary leptons l can couple to N via Yukawa terms analogous to (5), $y_{ia}l_iN_a\phi + $ h.c. However, for N being gauge singlets, the mirror leptons l' can couple with them with the same way, $y'_{ia}l'_iN_a\phi' + $ h.c.. In this way, the Ns play the role of messengers between ordinary and mirror particles. The whole set of relevant terms in the Lagrangian has the form

$$y_{ia}l_iN_a\phi + y'_{ia}l'_iN_a\phi' + \frac{M}{2}\,g_{ab}N_aN_b + \text{h.c.} \; . \qquad (19)$$

After integrating out the heavy states N, all operators (14), (15) and (16) are induced with the coupling constants

$$A = yg^{-1}y^T , \quad A' = y'g^{-1}y'^T , \quad D = yg^{-1}y'^T . \qquad (20)$$

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[i] Notice that the number of heavy singlet neutrinos, n, does not have to coincide with the number of standard families $n_g = 3$. From phenomenological constraints, $n = 2$ or $n > 3$ are also possible. In this view, generically g is a symmetric $n \times n$ matrix while y and y' are $3 \times n$ matrices.

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Without loss of generality, the matrix G can be taken real and diagonal. Then, assuming for simplicity that all states N_a have positive mirror parity, $MP : N_a \to N_a$, when $l_i \to \tilde{l}_i$, we see that MP requires that $y'_{ia} = y^*_{ia}$, and so for the coupling constants of the effective operators we obtain constraints (18).[j]

In the next sections we show that the N states can mediate L and CP violating scattering processes between the O- and M-sectors which could provide a new mechanism for primordial leptogenesis.

Either type of parity implies that the two sectors have the same particle physics.[k] If the two sectors are separate and do not interact by forces other than gravity, the difference between D and P parities is rather symbolic and does not have any profound implications. However, in scenarios with some particle messengers between the two sectors this difference can be important and can have dynamical consequences.

4. Spontaneously broken Mirror Parity

There is no reason to expect that Nature does not apply the old principle "The only good parity is a broken parity". It is interesting to discuss under which circumstances MP could be broken.

If O- and M-sectors are described by the minimal Standard Model, with the Higgs doublets ϕ and ϕ' having identical Higgs potentials

$$\mathcal{V}_{\text{Higgs}} = -\mu^2 \phi^\dagger \phi + \lambda^2 (\phi^\dagger \phi)^2 \,,$$
$$\mathcal{V}'_{\text{Higgs}} = -\mu^2 \phi'^\dagger \phi' + \lambda^2 (\phi'^\dagger \phi')^2 \tag{21}$$

then the VEVs of ϕ and ϕ' are unavoidably identical, $v = v' = \mu/\lambda$. In addition, from the experimental limits on the Higgs mass one can conclude that $\lambda \sim 1$.

The gauge symmetry of the theory allows also a quartic interaction term between O-and M-Higgs doublets,

$$\mathcal{V}_{\text{mix}} = \kappa (\phi^\dagger \phi)(\phi'^\dagger \phi') \,. \tag{22}$$

This term is cosmologically dangerous, since it would bring the two sectors into equilibrium in the early Universe via interactions $\bar{\phi}\phi \to \bar{\phi}'\phi'$ unless κ is

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[j] In general, some of N_a can have positive and others negative parity, i.e. $N_a \to p_a N_a$ where $p_a = \pm 1$. This would lead to $y'_{ia} = p_a y^*_{ia}$.

[k] The mirror parity could be spontaneously broken and the weak interaction scales $\langle \phi \rangle = v$ and $\langle \phi' \rangle = v'$ could be different, which leads to somewhat different particle physics in the mirror sector. The models with spontaneously broken parity and their implications were considered in Refs. [13, 14, 18]. In this paper we mostly concentrate on the case with exact mirror parity.

very small, $\kappa < 10^{-8}$ [14]. As far as $\kappa \ll \lambda$, this mixed term cannot cause the asymmetry between v and v'.

For achieving the breaking of mirror parity, one has to extend the Higgs sector. The simplest way is to introduce a real singlet scalar η which is odd under under mirror parity; it changes sign when the two gauge sectors exchange places, $\eta \to -\eta$. Therefore, its interaction Lagrangian with the O- and M-Higgs doublets should include terms $\propto \eta(\phi^\dagger \phi - \phi'^\dagger \phi')$. Therefore, if η has a non-zero VEV, it would induce a difference in mass-squared terms of ϕ and ϕ' and hence the mirror weak scale can be different from the ordinary one. But such an extension of the Higgs sector is not the most beautiful thing that one can imagine.

The situation remains the same if both O- and M-sectors are described by MSSM. In this case ordinary Higgses $\phi_{u,d}$ as well as their mirror partners $\phi'_{u,d}$ become chiral superfields and thus their renormalizable Lagrangian can include no mixed terms. Namely, the F-terms read as

$$\mathcal{L}_{\text{Higgs}} = \int d^2\theta(\mu\phi_u\phi_d + \mu\phi'_u\phi'_d) \ + \ \text{h.c.} \tag{23}$$

and clearly also gauge D-terms of O- and M-sectors are unmixed. In addition, experimental limits imply that μ-terms are of the order of 100 GeV.

The minimal gauge invariant term between the O- and M-Higgses in the superpotential has dimension 5: $(1/M)(\phi_u\phi_d)(\phi'_u\phi'_d)$, where M is some cutoff mass, e.g. of the order of the GUT or Planck scale. If this term is included together with (23), mixed quartic terms similar to (22) emerge in the Lagrangian,

$$\lambda(\phi_u^\dagger\phi_u)(\phi'_u\phi'_d) + \lambda(\phi_d^\dagger\phi_d)(\phi'_u\phi'_d) + (\phi_{u,d} \to \phi'_{u,d}) + \text{h.c.} \tag{24}$$

with the coupling constant $\lambda = \mu/M \ll 1$.

None of the soft supersymmetry breaking F-term and D-terms are dangerous. For example, the F-term $\frac{1}{M}\int d^2\theta z(\phi_u\phi_d)(\phi'_u\phi'_d)+$ h.c., where $z = m_S\theta^2$ being the supersymmetry breaking spurion with $m_S \sim 100$ GeV, gives rise to a quartic scalar term

$$\lambda(\phi_u\phi_d)(\phi'_u\phi'_d) + \text{h.c.} \tag{25}$$

with $\lambda \sim m_S/M \ll 1$. Thus for $\mu, m_S \sim 100$ GeV, all these quartic constants are strongly suppressed, and hence are safe.

It is easy, however, to construct a simple supersymmetric model where mirror parity is spontaneously broken. Let us introduce an additional singlet

superfield S and consider the Higgs superpotential having the form

$$W = \lambda S(\phi_u \phi_d + \phi'_u \phi'_d - \mathcal{M}^2) + MS^2 + ..., \tag{26}$$

where M and $\mathcal{M}$ are some large mass scales, much larger than M_W. Let us first consider the minimum of such a theory in the supersymmetric limit, without taking into account the supersymmetric breaking terms (soft masses and trilinear A-terms).

It is simple to see that such a theory can lead to the spontaneous breaking of M-parity due to different VEVs of the O- and M- Higgs doublets. Namely, mirror Higgses can get VEVs $v'_u = v'_d = \mathcal{M}$ while the ordinary ones get VEVs of the order of the supersymmetry breaking scale. The hierarchy problem between two VEVs is solved via the so called Pseudo-Goldstone boson mechanism [23].

Then the F-term condition $F_S = 0$ shows that $v_u v_d + v'_u v'_d = \mathcal{M}^2$, but does not fix how the VEVs are distributed between $\phi_{u,d}$ and $\phi'_{u,d}$ Higgses. In fact, this superpotential has an accidental global symmetry $SU(4)$, larger than the local symmetry $SU(2) \times SU(2)'$. Two "upper" doublets (ϕ_u, ϕ'_u) "down" doublets as (ϕ_d, ϕ'_d), form 4-plets of $SU(4)$. So if one chooses that $v_{u,d} = 0$, than we have $v'_u v'_d = \mathcal{M}^2$, and in addition from the D-term condition of $SU(2)'$ it follows that $v'_u = v'_d = v'/\sqrt{2}$, i.e. $\tan \beta' = 1$. As for the doublets $\phi_{u,d}$, in the limit of unbroken supersymmetry they remain as massless Goldstone modes of the spontaneously broken accidental symmetry $SU(4)$, and their masses/VEVs can emerge only after the supersymmetry breaking. In addition, a non-zero μ-term is generated since S will get a VEV of the order of m_S.

In the Standard Model, with one Higgs ϕ, the electroweak breaking scale identified with the Higgs VEV is unambiguously fixed by the Fermi constant, $v = 174$ GeV, and we have to take this into account seriously. In the models with two Higgs doublets, ϕ_u and ϕ_d, the VEVs v_u and v_d cannot be known separately, the parameter $\tan \beta = v_u/v_d$ can be arbitrary (well, with all probability larger than 1 but smaller than 100), however the total VEV $v_u^2 + v_d^2 = v^2$.

Let us discuss now what happens if MP is spontaneously broken at lower energies by two electroweak scales $\langle \phi' \rangle = v'$ and $\langle \phi \rangle = v$ to be different; we assume that $v' \gg v$. As far as the Yukawa couplings have the same values in both systems, the mass and mixing pattern of the charged fermions in the mirror world is completely analogous to that of the visible one, but with all fermion masses scaled up by the factor $\zeta = v'/v$. The masses of gauge bosons and Higgses are also scaled as $M_{W',Z',\phi'} = \zeta M_{W,Z,\phi}$ while photons

and gluons remain massless in both sectors.

With regard to the two chromodynamics, a big difference between the electroweak scales v' and v will not cause a similar big difference between the confinement scales in two worlds. Indeed, if P parity is valid at higher (GUT) scales, the strong coupling constants in both sectors would evolve down in energy with the same values until the energy reaches the value of the mirror-top (t') mass. Below that α'_s will have a different slope from α_s. It is then very easy to calculate the value of the scale Λ' at which α'_s becomes large. This value of course depends on the ratio $\zeta = v'/v$.

Taking $\Lambda = 200$ MeV for the ordinary QCD, we find $\Lambda' \simeq 280$ MeV if $\zeta \sim 30$. On the other hand, we have $m'_{u,d} = \zeta m_{u,d} \sim m_s$ so that masses of the mirror light quarks u' and d' do not exceed Λ'. So the condensates $\langle \bar{q}' q' \rangle$ should be formed with approximately the same magnitudes as the usual quark condensates $\langle \bar{q} q \rangle$. As a result, mirror pions should have mass $m'_\pi \simeq \sqrt{m_{u',d'} \langle \bar{q}' q' \rangle}$ comparable to the mass of normal kaons $m_K \simeq \sqrt{m_s \langle \bar{q} q \rangle}$.

As for the mirror nucleons, their masses are approximately 1.5 times larger than that of the usual nucleons. Since $(m'_d - m'_u) \approx 30(m_d - m_u)$ we expect the mirror neutron n' to be heavier than the mirror proton p' by about 150 MeV or so, while the mirror electron mass is $m'_e = \zeta m_e \sim 15$ MeV. Clearly, such a large mass difference cannot be compensated by the nuclear binding energy and hence even bound neutrons will be unstable against β decay $n' \to p' e' \bar{\nu}'_e$. Thus in the mirror world hydrogen will be the only stable nucleus.

Therefore, the microphysics of the mirror world would not be the same as that of the ordinary world, and the M-sector essentially becomes a sort of hidden sector with particle properties which only can be guessed (deduced) to some extent from the coupling constant structures of the ordinary world. It is also a good exercise to understand what could happen to our world should the electroweak scale be different.

5. Kinetic mixing of Ordinary and Mirror Photons

In the context of $G_{\mathrm{SM}} \times G'_{\mathrm{SM}}$, the Lagrangian can contain the gauge invariant mixing term between the field-strength tensors of the gauge factors $U(1)$ and $U(1)'$. After the electroweak symmetry breaking, this term gives rise to a kinetic mixing term between the the O- and M-photons,

$$\mathcal{L} = -\varepsilon F^{\mu\nu} F'_{\mu\nu}\,. \tag{27}$$

This term cannot be suppressed for symmetry reasons, and generally the constant ε could be of order 1.
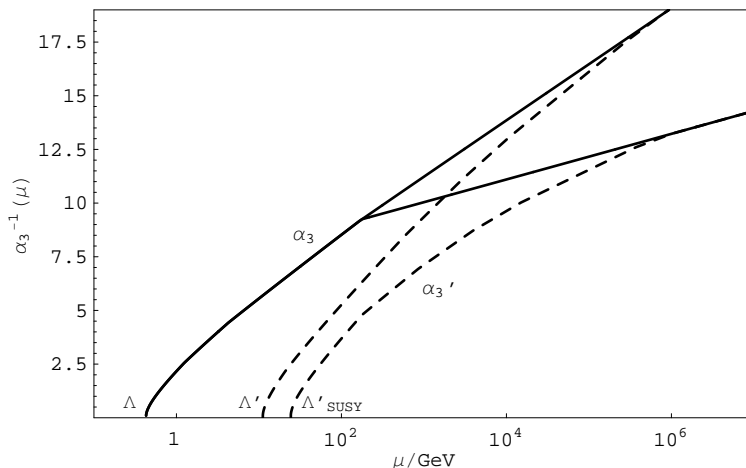
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Figure 1.   Evolution of the strong coupling constant in the two sectors from high energies to lower energies in the case of $v' = 10^6$ GeV for the Standard Model case and the case of MSSM with Supersymmetry broken at 200 GEV (upper an lower respectively). The solid curve stands for the O-sector and the dashed curve for the M-sector.

Once such a term is introduced, the following situation emerges. One has to diagonalize first the kinetic terms of the ordinary photon field A_μ and the mirror one A'_μ, and identify the physical photon as a linear combination. Now, once the kinetic terms are brought to canonical form by diagonalization and scaling of the fields, $(A, A') \to (A_1, A_2)$, any orthonormal combination of states A_1 and A_2 becomes good to describe the physical basis. In particular, A_2 can be chosen as a "sterile" state which does not couple to O-particles but only to M-particles. Then the orthogonal combination A_1 couples not only to O-particles, but also with M-particles with a small charge $\propto 2\varepsilon$ – in other words, mirror matter becomes "millicharged" with respect to the physical ordinary photon [7, 21].

In this way, the term (27) induces the process $e^+e^- \to e'^+e'^-$ with an amplitude just 2ε times the s-channel amplitude for $e^+e^- \to e^+e^-$. This could have striking experimental implications for positronium physics; ordinary positronium mixes to its mirror counterpart which effect could be manifested as an invisible decay mode of the orthopositronium. Perhaps this effect could be important for the troubling mismatch problems in orthopositronium physics [8, 24, 25]. For the moment, the experimental limits on the orthopositronium decays lead to an upper limit $\varepsilon < 3 \times 10^{-7}$ or so.

A stronger limit can be obtained from the cosmology. As we already remarked, the BBN constraints require that the mirror sector should be colder than the the ordinary one, $T'/T < 0.5$ or so. On the other hand, the

reaction $e^+e^- \to e'^+e'^-$, funneling energy from the O-sector to the M-sector, would heat the latter too much before the BBN epoch, unless $\varepsilon < 3 \times 10^{-8}$ [9].

The search for the process $e^+e^- \to invisible$ could approach sensitivities down to few $\times 10^{-9}$. [26] This interesting experiment could test the proposal of Ref. [27] claiming that the signal for dark matter detection by the DAMA/NaI group [28] can be explained by elastic scattering of M-baryons with ordinary ones mediated by kinetic mixing (27), if $\varepsilon \sim 4 \times 10^{-9}$.

The smallness of the kinetic mixing term (27) can be naturally explained by invoking the concept of grand unification. Obviously, the term (27) is forbidden in GUTs like $SU(5) \times SU(5)'$ which do not contain abelian factors. However, given that both $SU(5)$ and $SU(5)'$ symmetries are broken down to their $SU(3) \times SU(2) \times U(1)$ subgroups by the Higgs 24-plets Φ and Φ', it could emerge from the higher order effective operator

$$\mathcal{L} = -\frac{\zeta}{M^2} (G^{\mu\nu}\Phi)(G'_{\mu\nu}\Phi') \,, \tag{28}$$

where $G_{\mu\nu}$ and $G'_{\mu\nu}$ are field-strength tensors respectively of $SU(5)$ and $SU(5)'$, and M is some cutoff scale which can be of the order of M_{Pl} or so. After substituting VEVs of Φ and Φ' the operator (27) is induced with $\varepsilon \sim \zeta(\langle\Phi\rangle/M)^2$.

In fact, the operator (28) can be effectively induced by loop-effects involving some heavy fermion or scalar fields in the mixed representations of $SU(5) \times SU(5)'$, with $\zeta \sim \alpha/3\pi \sim 10^{-3}$ being a loop-factor. Consider, for example, $SU(5) \times SU(5)'$ theory which apart from the standard O- and M-fermion multiplets includes also the chiral fermions in mixed representations $F \sim (5,5)$ and $F' \sim (\bar{5},\bar{5})$. These would necessarily appear if $SU(5) \times SU(5)'$ is embedded into e.g. the $SU(10)$ group. They should have a large mass term MFF', e.g. of the order of the $SU(10)$ breaking scale to $SU(5) \times SU(5)'$. However, in general they could have coupling terms $\Phi FF' + \Phi'FF'$ with the GUT Higgses $\Phi \sim (24,1)$ and $\Phi' \sim (1,24)$.

With respect to the $G_{\mathrm{SM}} \times G'_{\mathrm{SM}}$ subgroup, these multiplets split into fragments F_{ij} with different hypercharges (Y_i, Y'_j) with respect to $U(1)$ and $U(1)'$ factors, and correspondingly with masses $\hat{M}_{ij} = M + Y_i\langle\Phi\rangle + Y'_j\langle\Phi'\rangle$. Therefore, the loops involving the fermions F_{ij} would induce a contribution to the term (27) with $\varepsilon \simeq (\alpha/3\pi)\mathrm{Tr}[YY'\ln(\hat{M}/\Lambda)]$ where Λ is an ultraviolet cutoff scale and in the trace a sum over all fragments F_{ij} is understood. As far as these fragments emerge from the GUT multiplets, they necessarily obey that $\mathrm{Tr}(YY') = 0$, and thus ε should be finite and cutoff independent. Thus, expanding the logarithm in terms of small parameters $\langle\Phi^{(\prime)}\rangle/M$, we

finally obtain

$$\varepsilon \simeq \frac{\alpha \langle \Phi \rangle \langle \Phi' \rangle}{3\pi M^2} \, \mathrm{Tr} \left[(YY')^2 \right], \tag{29}$$

exactly what we expected from the effective operator (28). Hence, the heavy mixed multiplets in fact do not decouple and induce the O- and M-photon kinetic mixing term proportional to the square of typical mass splittings in these multiplets ($\sim \langle \Phi \rangle^2$), analogously to the familiar situation for the photon to Z-boson mixing in the standard model. Hence, taking the GUT scale as $\langle \Phi \rangle \sim 10^{16}$ GeV and $M \sim M_{Pl}$ we see that the strength of the kinetic mixing term (27) could vary vary from $\varepsilon \sim 10^{-10}$ to 10^{-8}. Certainly, the coupling (28) can be stronger suppressed or completely eliminated by some symmetry reasons.

6. Other possible interactions between O- and M-particles

Here we briefly discuss what other common interactions and forces could exist between the O- and M-particles, including matter fields and gauge fields.

It is quite possible that O-and M-particles have common forces mediated by the gauge bosons of some additional symmetry group H. In other words, one can consider a theory with a gauge group $G \times G' \times H$, where O-particles are in some representations of H, $L_a \sim r_a$, and correspondingly their antiparticles are in antirepresentations, $\tilde{R}_a \sim \bar{r}_a$. As for M-particles, vice versa, we take $L'_a \sim \bar{r}_a$, and so $\tilde{R}'_a \sim r_a$. Only such a prescription of $\mathcal{G}$ pattern is compatible with the mirror parity (9). In addition, in this case H symmetry automatically becomes vector-like and so it would have no problems with axial anomalies even if the particle contents of O- and M-sectors separately are not anomaly-free with respect to H.

Let us consider the following example. The horizontal flavor symmetry $SU(3)_H$ between the quark-lepton families seems to be very promising for understanding the fermion mass and mixing pattern [30, 31]. In addition, it can be useful for controlling the flavor-changing phenomena in the context of supersymmetry [15]. One can consider e.g. a GUT with $SU(5) \times SU(3)_H$ symmetry where L-fermions in (11) are triplets of $SU(3)_H$. So $SU(3)_H$ has a chiral character and it is not anomaly-free unless some extra states are introduced for the anomaly cancellation [30].

However, the concept of mirror sector makes the things easier. Consider e.g. the $SU(5) \times SU(5)' \times SU(3)_H$ theory with L-fermions being triplets of $SU(3)_H$ and L'-fermions anti-triplets. Hence, in this case the $SU(3)_H$

anomalies of the ordinary particles are cancelled by their mirror partners. Another advantage is that in a supersymmetric theory the gauge D-terms of $SU(3)_H$ are perfectly cancelled between the two sectors and hence they do not give rise to dangerous flavor-changing phenomena [15]. Common gauge $B - L$ symmetry between the two sectors is also plausible.

The immediate implication of interactions mediated by common gauge or Higgs bosons would be the mixing of neutral O-bosons to their M-partners, mediated by horizontal gauge bosons. Namely, oscillations $\pi^0 \to \pi'^0$ or $K^0 \to K'^0$ become possible and perhaps even detectable if the horizontal symmetry breaking scale is not too high.

The operators of dimension 9, $(1/M^5)(udd)(udd)$, would lead to oscillations between ordinary and mirror neutrons, $n \to n'$. Obviously, the phenomenological limits on such oscillations are comparable to that of neutron-antineutron oscillation which is restricted by the limits on the nucleon stability.

The model with common Peccei–Quinn symmetry between the O- and M-sectors was considered in [16]. In this situation many astrophysical and cosmological bounds on the axion can be eliminated.

7. The expansion of the Universe and thermodynamics of the O- and M-sectors

Let us assume that, after inflation ended, the O- and M-systems received different reheating temperatures, namely $T_R > T_R'$. This is certainly possible despite the fact that the two sectors have identical Lagrangians, and can be naturally achieved in certain models of inflation [14, 17, 18].[1]

If the two systems were decoupled already after reheating, at later times t they will have different temperatures $T(t)$ and $T'(t)$, and so different energy and entropy densities:

$$\rho(t) = \frac{\pi^2}{30} g_*(T)T^4, \quad \rho'(t) = \frac{\pi^2}{30} g_*'(T')T'^4 , \tag{30}$$

$$s(t) = \frac{2\pi^2}{45} g_s(T)T^3, \quad s'(t) = \frac{2\pi^2}{45} g_s'(T')T'^3 . \tag{31}$$

The factors g_*, g_s and g_*', g_s' accounting for the effective number of the degrees of freedom in the two systems can in general be different from each other. Let us assume that during the expansion of the Universe the two

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[1] For analogy, two harmonic oscillators with the same frequency (e.g. two springs with the same material and the same length) are not obliged to oscillate with the same amplitudes.

sectors evolve with separately conserved entropies. Then the ratio $x \equiv (s'/s)^{1/3}$ is time independent while the ratio of the temperatures in the two sectors is simply given by

$$\frac{T'(t)}{T(t)} = x \cdot \left[\frac{g_s(T)}{g'_s(T')}\right]^{1/3} . \tag{32}$$

The Hubble expansion rate is determined by the total energy density $\bar{\rho} = \rho + \rho'$, $H = \sqrt{(8\pi/3)G_N\bar{\rho}}$. Therefore, at a given time t in a radiation dominated epoch we have

$$H(t) = \frac{1}{2t} = 1.66\sqrt{\bar{g}_*(T)}\,\frac{T^2}{M_{Pl}} = 1.66\sqrt{\bar{g}'_*(T')}\,\frac{T'^2}{M_{Pl}} \tag{33}$$

in terms of O- and M-temperatures $T(t)$ and $T'(t)$, where

$$\bar{g}_*(T) = g_*(T)(1 + x^4), \qquad \bar{g}'_*(T') = g'_*(T')(1 + x^{-4}). \tag{34}$$

In particular, we have $x = T'_0/T_0$, where T_0, T'_0 are the present temperatures of the relic photons in O- and M-sectors. In fact, x is the only free parameter in our model and it is constrained by the BBN bounds.

The observed abundances of light elements are in good agreement with the standard nucleosynthesis predictions. At $T \sim 1$ MeV we have $g_* = 10.75$ as it is saturated by photons γ, electrons e and three neutrino species $\nu_{e,\mu,\tau}$. The contribution of mirror particles (γ', e' and $\nu'_{e,\mu,\tau}$) would change it to $\bar{g}_* = g_*(1 + x^4)$. Deviations from $g_* = 10.75$ are usually parametrized in terms of the effective number of extra neutrino species, $\Delta g = \bar{g}_* - 10.75 = 1.75 \cdot \Delta N_\nu$. Thus we have

$$\Delta N_\nu = 6.14 \cdot x^4 . \tag{35}$$

This limit very weakly depends on ΔN_ν. Namely, the conservative bound $\Delta N_\nu < 1$ implies $x < 0.64$. In view of the present observational situation, confronting the WMAP results with the BBN analysis, the bound seems to be stronger. However, e.g. $x = 0.3$ implies a completely negligible contribution $\Delta N_\nu = 0.05$.

As far as $x^4 \ll 1$, in a relativistic epoch the Hubble expansion rate (33) is dominated by the O-matter density and the presence of the M-sector practically does not affect the standard cosmology of the early ordinary Universe. However, even if the two sectors have the same microphysics, the cosmology of the early mirror world can be very different from the standard one as far as the crucial epochs like baryogenesis, nucleosynthesis, etc. are concerned. Any of these epochs is related to an instant when the rate of the relevant particle process $\Gamma(T)$, which is generically a function of the

temperature, becomes equal to the Hubble expansion rate $H(T)$. Obviously, in the M-sector these events take place earlier than in the O-sector, and as a rule, the relevant processes in the former freeze out at larger temperatures than in the latter.

In the matter domination epoch the situation becomes different. In particular, we know that ordinary baryons provide only a small fraction of the present matter density, whereas the observational data indicate the presence of dark matter with about 5 times larger density. It is interesting to question whether the missing matter density of the Universe could be due to mirror baryons? In the next section we show that this could occur in a pretty natural manner.

It can also be shown that the BBN epoch in the mirror world proceeds differently from the ordinary one, and it predicts different abundances of primordial elements [19]. It is well known that primordial abundances of the light elements depend on the baryon to photon density ratio $\eta = n_B/n_\gamma$, and the observational data well agree with the WMAP result $\eta \simeq 6 \times 10^{-10}$. As far as $T' \ll T$, the universe expansion rate at the ordinary BBN epoch ($T \sim 1$ MeV) is determined by the O-matter density itself, and thus for the ordinary observer *Olga* it would be very difficult to detect the contribution of the M-sector; the latter is equivalent to $\Delta N_\nu \approx 6.14 x^4$ and hence it is negligible for $x \ll 1$. As for nucleosynthesis epoch in M-sector, the contribution of O-world instead is dramatic: it is equivalent to $\Delta N'_\nu \approx 6.14 x^{-4} \gg 1$. Therefore, the mirror observer *Maxim* who measures the abundances of mirror light elements should immediately observe a discrepancy between the universe expansion rate and the M-matter density at his BBN epoch ($T' \sim 1$ MeV) as far as the former is determined by O-matter density which is invisible for *Maxim*. The result for mirror ^{4}He also depends on the mirror baryon to photon density ratio $\eta' = n'_B/n'_\gamma$. Recalling that $\eta' = (\beta/x^3)\eta$, we see that $\eta \gg \eta$ unless $\beta = n'_B/n_B \ll x^3$. However, if $\beta > 1$, we expect that the mirror helium mass fraction Y'_4 would be considerably larger than the observable $Y_4 \simeq 0.24$. Namely, direct calculations show that for x varying from 0.6 to 0.1, Y'_4 would vary in the range $Y'_4 = 0.5 - 0.8$. Therefore, if M-baryons constitute dark matter or at least its reasonable fraction, the M-world is a dominantly helium world while the heavier elements can also present with significant abundances.
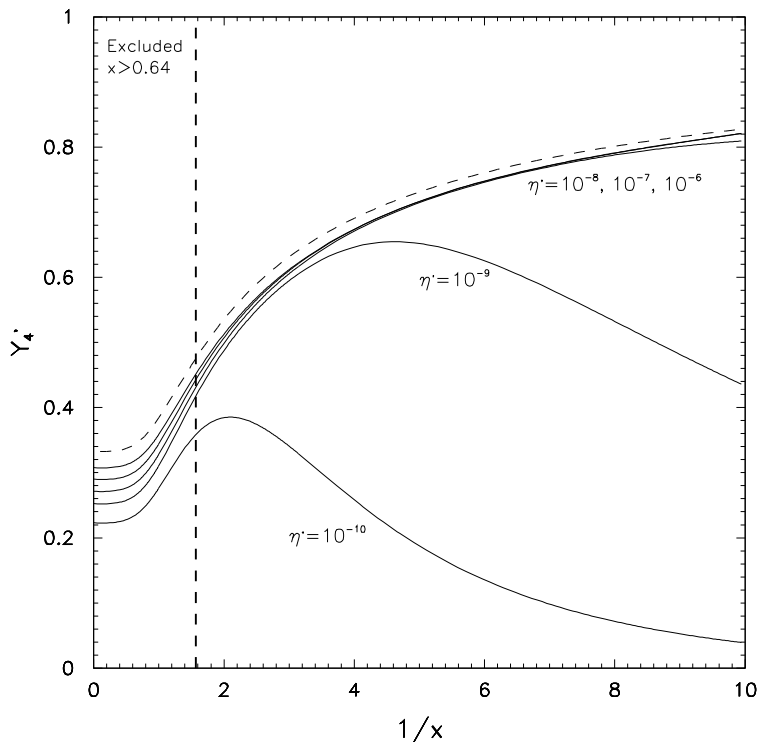
Figure 2.   The primordial mirror ^{4}He mass fraction as a function of x. The different curves correspond to different fixed values of of $\eta' = n'_B/n'_\gamma$. The dashed curve extrapolates the case when mirror baryons constitute dark matter, i.e. $\beta \simeq 5$.

8. Baryogenesis in M-sector and mirror baryons as dark matter

8.1. *Visible and dark matter in the Universe*

The present cosmological observations strongly support the main predictions of the inflationary scenario: first, the Universe is flat, with the energy density very close to the critical $\Omega = 1$, and second, primordial density perturbations have a nearly flat spectrum, with the spectral index $n_s \approx 1$. The non-relativistic matter gives only a small fraction of the present energy density, about $\Omega_m \simeq 0.27$, while the rest is attributed to the vacuum energy (cosmological term) $\Omega_\Lambda \simeq 0.73$ [32]. The fact that Ω_m and Ω_Λ are of the same order, gives rise to so called cosmological coincidence problem: why do we live in an epoch where $\rho_m \sim \rho_\lambda$, if in the early Universe one had $\rho_m \gg \rho_\Lambda$ and in the late Universe one would expect $\rho_m \ll \rho_\Lambda$? The answer can be only related to an anthrophic principle: the matter and vacuum energy densities scale differently with the expansion of the Universe $\rho_m \propto a^{-3}$ and $\rho_\Lambda \propto$

const., hence they have to coincide at some moment, and we are just happy to be here. Moreover, for substantially larger ρ_Λ no galaxies could be formed and thus there would not be anyone to ask this this question.

On the other hand, the matter in the Universe has two components, visible and dark, $\Omega_m = \Omega_b + \Omega_d$. The visible matter consists of baryons with $\Omega_b \simeq 0.044$ while the dark matter with $\Omega_d \simeq 0.22$ is constituted by some hypothetic particle species very weakly interacting with the observable matter. It is a tantalizing question, why do the visible and dark components have such close energy densities? Clearly, the ratio

$$\beta = \frac{\rho_d}{\rho_b} \tag{36}$$

does not depend on time as far as with the expansion of the Universe both ρ_b and ρ_d scale as $\propto a^{-3}$.

In view of the standard cosmological paradigm, there is no good reason for having $\Omega_d \sim \Omega_b$, as the visible and dark components have different origins. The density of the visible matter is $\rho_b = M_N n_b$, where $M_N \simeq 1$ GeV is the nucleon mass, and n_b is the baryon number density of the Universe. The latter should be produced in a very early Universe by some baryogenesis mechanism, which is presumably related to some B and CP-violating physics at very high energies. The baryon number per photon $\eta = n_b/n_\gamma$ is very small. Observational data on the primordial abundances of light elements and the WMAP results on the CMBR anisotropies nicely converge to the value $\eta \approx 6 \times 10^{-10}$.

As for dark matter, it is presumably constituted by some cold relics with mass M and number density n_d, and $\rho_d = M n_d$. The most popular candidate for cold dark matter (CDM) is provided by the lightest supersymmetric particle (LSP) with $M_{\mathrm{LSP}} \sim 1$ TeV, and its number density n_{LSP} is fixed by its annihilation cross-section. Hence $\rho_b \sim \rho_{\mathrm{LSP}}$ requires that $n_b/n_{\mathrm{LSP}} \sim M_{\mathrm{LSP}}/M_N$ and the origin of such a conspiracy between four principally independent parameters is absolutely unclear. Once again, the value M_N is fixed by the QCD scale while M_{LSP} is related to the supersymmetry breaking scale, n_b is determined by B and CP violating properties of the particle theory at very high energies whereas n_{LSP} strongly depends on the supersymmetry breaking details. Within the parameter space of the MSSM it could vary within several orders of magnitude, and moreover, in either case it has nothing to do with the B and CP violating effects.

The situation looks even more obscure if the dark component is related e.g. to the primordial oscillations of a classical axion field, in which case the dark matter particles constituted by axions are superlight, with

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mass $\ll 1$ eV, but they have a condensate with enormously high number density.

In this view, the concept of mirror world could give a new twist to this problem. Once the visible matter is built up by ordinary baryons, then the mirror baryons could constitute dark matter in a natural way. They interact with mirror photons, however they are dark in terms of the ordinary photons. The mass of M-baryons is the same as the ordinary one, $M = M_N$, and so we have $\beta = n'_b/n_b$, where n'_b is the number density of M-baryons. In addition, as far as the two sectors have the same particle physics, it is natural to think that the M-baryon number density n'_b is determined by the baryogenesis mechanism which is similar to the one which fixes the O-baryon density n_b. Thus, one could question whether the ratio $\beta = n'_b/n_b$ could be naturally order 1 or somewhat bigger.

The visible matter in the Universe consists of baryons, while the abundance of antibaryons is vanishingly small. In the early Universe, at temperatures $T \gg 1$ GeV, the baryons and antibaryons had practically the same densities, $n_b \approx n_{\bar{b}}$ with n_b slightly exceeding $n_{\bar{b}}$, so that the baryon number density was small, $n_B = n_b - n_{\bar{b}} \ll n_b$. If there was no significant entropy production after the baryogenesis epoch, the baryon number density to entropy density ratio had to be the same as today, $B = n_B/s \approx 8 \times 10^{-11}$.[m]

One can question, by whom, and how, was the initial Universe prepared with such a small excess of baryons over antibaryons. In the Friedman Universe the initial baryon asymmetry could be arranged a priori, in terms of non-vanishing chemical potential of baryons. However, the inflationary paradigm gives another twist to this question, since inflation dilutes any pre-existing baryon number of the Universe to zero. Therefore, after inflaton decay and the (re-)heating of the Universe, the baryon asymmetry has to be created by some cosmological mechanism.

There are several relatively honest baryogenesis mechanisms such as GUT baryogenesis, leptogenesis, electroweak baryogenesis, etc. (for a review, see e.g. [33]). They are all based on general principles suggested a long time ago by Sakharov [34]. A non-zero baryon asymmetry can be produced in the initially baryon symmetric Universe if three conditions are fulfilled: B-violation, C- and CP-violation and departure from thermal equilibrium. In the GUT baryogenesis or leptogenesis scenarios these conditions can be satisfied in the decays of heavy particles.

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[m] In the following we use $B = n_B/s$ which is related with the familiar $\eta = n_B/n_\gamma$ as $B \approx 0.14\eta$. However, B is better suited to characterizing the baryon asymmetry since it does not depend on time if the entropy of the Universe is conserved.

At present it is not possible to say definitely which of the known mechanisms is responsible for the observed baryon asymmetry in the ordinary world. However, it is most likely that the baryon asymmetry in the mirror world is produced by the same mechanism and, moreover, the properties of the B and CP violation processes are parametrically the same in both cases. But the mirror sector has a lower temperature than the ordinary one, and so at epochs relevant for baryogenesis the out-of-equilibrium conditions should be more easily fulfilled for the M-sector.

8.2. *Baryogenesis in the O- and M-worlds*

Let us consider the difference between the ordinary and mirror baryon asymmetries on the example of the GUT baryogenesis mechanism. It is typically based on 'slow' B- and CP-violating decays of a superheavy boson X into quarks and leptons, where slow means that at $T < M$ the Hubble parameter $H(T)$ is greater than the decay rate $\Gamma \sim \alpha M$, α being the coupling strength of X to fermions and M its mass. The other reaction rates are also of relevance: *inverse decay*, $\Gamma_I \sim \Gamma (M/T)^{3/2} \exp(-M/T)$ for $T < M_X$; and *the X boson mediated scattering processes*, $\Gamma_S \sim n_X \sigma \sim A\alpha^2 T^5/M^4$, where the factor A accounts for the possible reaction channels.

The final BA depends on the temperature at which X bosons go out of equilibrium. One can introduce a parameter which measures the effectiveness of the decay at the epoch $T \sim M$, $k = (\Gamma/H)_{T=M} = 0.3\bar{g}_*^{-1/2}(\alpha M_{Pl}/M)$. For $k \ll 1$ the out-of-equilibrium condition is strongly satisfied, and per decay of one X particle one generates the baryon number proportional to the CP-violating asymmetry ε. Thus, we have $B = \varepsilon/g_*$, g_* is a number of effective degrees of freedom at $T < M$. The larger k is, the longer equilibrium is maintained and the freeze-out abundance of the X boson becomes smaller. Hence, the resulting baryon number to entropy ratio becomes $B = (\varepsilon/g_*)D(k)$, where the damping factor $D(k)$ is a decreasing function of k. In particular, $D(k) = 1$ for $k \ll 1$, while for k exceeding some critical value k_c, the damping is exponential.

The presence of the mirror sector practically does not alter the ordinary baryogenesis. The effective particle number is $\bar{g}_*(T) = g_*(T)(1 + x^4)$ and thus the contribution of M-particles to the Hubble constant at $T \sim M$ is suppressed by a small factor x^4.

In the mirror sector everything should occur in a similar way, apart from the fact that now at $T' \sim M$ the Hubble constant is not dominated by the mirror species but by ordinary ones, $\bar{g}'_*(T') \simeq g'_*(T')(1 + x^{-4})$. As a consequence, we have $k' = (\Gamma/H)_{|T'=M} = kx^2$. Therefore, the damping factor for

mirror baryon asymmetry can be simply obtained by replacing $k \to k' = kx^2$ in $D(k)$. In other words, the baryon number density to entropy density ratio in the M-world becomes $B' = n'_B/s' \simeq (\epsilon/g_*)D(kx^2)$. Since $D(k)$ is a decreasing function of k, then for $x < 1$ we have $D(kx^2) > D(k)$ and thus we conclude that the mirror world always gets a *larger* baryon asymmetry than the visible one, $B' > B$.[n] Namely, for $k > 1$ the baryon asymmetry in the O-sector is damped by some factor – we have $B \simeq (\varepsilon/g_*)D(k) < \varepsilon/g_*$, while if $x^2 < k^{-1}$, the damping would be irrelevant for the M-sector and hence $B' \simeq \varepsilon/g_*$.

However, this does not a priori mean that Ω'_b will be larger than Ω_b. Since the entropy densities are related as $s'/s = x^3$, for the ratio $\beta = \Omega'_b/\Omega_b$ we have

$$\beta(x) = \frac{n'_B}{n_B} = \frac{B's'}{Bs} = \frac{x^3 D(kx^2)}{D(k)} \ . \tag{37}$$

The behavior of this ratio as a function of k for different values of the parameter x is given in Ref. [19]. Clearly, in order to have $\Omega'_b > \Omega_b$, the function $D(k)$ has to decrease faster than $k^{-3/2}$ between $k' = kx^2$ and k. Closer inspection of this function reveals that the M-baryons can be overproduced only if k is sufficiently large, so that the relevant interactions in the observable sector maintain equilibrium longer than in the mirror one, and thus ordinary BA can be suppressed by an exponential Boltzmann factor while the mirror BA could be produced still in the regime $k' = kx^2 \ll 1$, when $D(k') \approx 1$.

However, the GUT baryogenesis picture has the following generic problem. In scenarios based on grand unification models like $SU(5)$, the heavy gauge or Higgs boson decays separately violate B and L, but conserve $B-L$, and so finally $B-L = 0$. On the other hand, the non-perturbative sphaleron processes, which violate $B + L$ but conserve $B - L$, are effective at temperatures from about 10^{12} GeV down to 100 GeV [35]. Therefore, if $B + L$ is erased by sphaleron transitions, the final B and L will both vanish.

Hence, in a realistic scenario one actually has to produce a non-zero $B - L$ rather than just a non-zero B, a fact that strongly favors the so called *leptogenesis* scenario [36]. The seesaw mechanism for neutrino masses offers the elegant possibility of generating non-zero $B - L$ in CP-violating decays of heavy Majorana neutrinos N into leptons and Higgses. These decays

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[n] As was shown in Ref. [19], the relation $B' > B$ appears also in the context of the electroweak baryogenesis scenario, where the out-of-equilibrium conditions are provided by fast phase transition and bubble nucleation.

violate L but obviously do not change B and so they could create a non-zero $B - L = -L_{\text{in}}$. Namely, due to complex Yukawa constants, the decay rates $\Gamma(N \to l\phi)$ and $\Gamma(N \to \tilde{l}\tilde{\phi})$ can be different from each other, so that the leptons l and anti-leptons $\tilde{l}$ are produced in different amounts.

When sphalerons are in equilibrium, they violate $B+L$ and so redistribute non-zero $B-L$ between the baryon and lepton numbers of the Universe. The final values of B and $B - L$ are related by $B = a(B - L)$, where a is order 1 coefficient, namely $a \simeq 1/3$ in the SM and in its supersymmetric extension [33]. Hence, the observed baryon to entropy density ratio, $B \approx 8 \times 10^{-11}$, needs to produce $B - L \sim 2 \times 10^{-10}$.

However, the comparative analysis presented above for the GUT baryogenesis in the O- and M-worlds, is essentially true also for the leptogenesis scenario. The out-of-equilibrium decays of heavy N neutrinos of the O-sector would produce a non-zero $B-L$ which being reprocessed by sphalerons would give an observable baryon asymmetry $B = a(B - L)$. On the other hand, the same decays of heavy N' neutrinos of the M-sector would give non-zero $(B' - L')$ and thus the mirror baryon asymmetry $B' = a(B' - L')$. In order to thermally produce heavy neutrinos in both O- and M-sectors, the lightest of them should have a mass smaller than the reheating temperature T'_R in the M-sector, i.e. $M_N < T'_R, T_R$. The situation $M_N > T'_R$ would prevent thermal production of N' states, and so no $B' - L'$ would be generated in the M-sector. However, one can also consider scenarios where both N and N' states are non-thermally produced in inflaton decays, but in different amounts. Then the reheating of both sectors as well as $B - L$ number generation can be related to the decays of the heavy neutrinos of both sectors and hence the situation $T'_R < T_R$ can be naturally accompanied by $B' > B$.

> *Tweedledum and Tweedledee*
> *Agreed to have a battle;*
> *For Tweedledum said Tweedledee*
> *Had spoiled his nice new rattle.*

9. Baryogenesis via Ordinary-Mirror particle interaction

An alternative mechanism of leptogenesis is based on scattering processes rather than on decay [20]. The main idea consists in the following. The hidden (mirror) sector of particles is not in thermal equilibrium with the ordinary particle world as far as the two systems interact very weakly. However, superheavy singlet neutrinos can mediate very weak effective interactions between the ordinary and mirror leptons. Then a net $B - L$ could emerge in the Universe as a result of CP-violating effects in the unbalanced production

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of mirror particles from ordinary particle collisions.

As far as the leptogenesis is concerned, we concentrate only on the lepton sector of both O and M worlds. Therefore we consider the Standard Model, and among other particles species, concentrate on the lepton doublets $l_i = (\nu, e)_i$ ($i = 1, 2, 3$ is the family index) and the Higgs doublet ϕ for the O-sector, and on their mirror partners $l'_i = (\nu', e')_i$ and ϕ'. Their couplings to the heavy singlet neutrinos are given by (19).

Let us now discuss this mechanism in more detail. A crucial role in our considerations is played by the reheating temperature T_R, at which the inflaton decay and entropy production of the Universe is over, and after which the Universe is dominated by a relativistic plasma of ordinary particle species. As we discussed above, we assume that after the postinflationary reheating, different temperatures are established in the two sectors: $T'_R < T_R$, i.e. the mirror sector is cooler than the visible one, or ultimately, even completely "empty".º

In addition, the two particle systems should interact very weakly so that they do not come into thermal equilibrium with each other after reheating. We assume that the heavy neutrino masses are larger than T_R and thus cannot be thermally produced. As a result, the usual leptogenesis mechanism via $N \to l\phi$ decays is ineffective.

Now, the important role is played by lepton number violating scatterings mediated by the heavy neutrinos N. The "cooler" mirror world starts to be "slowly" occupied due to entropy transfer from the ordinary sector through the $\Delta L = 1$ reactions $l_i\phi \to \bar{l}'_k\bar{\phi}'$, $\bar{l}_i\bar{\phi} \to l'_k\phi'$. In general these processes violate CP due to complex Yukawa couplings in Eq. (19), and so the cross-sections with leptons and anti-leptons in the initial state are different from each other. As a result, leptons leak to the mirror sector with different rate than antileptons and so a non-zero $B - L$ is produced in the Universe.

It is important to stress that this mechanism would generate the baryon asymmetry not only in the observable sector, but also in the mirror sector. In fact, the two sectors are completely similar, and have similar CP-violating properties. We have scattering processes which transform the ordinary particles into their mirror partners, and CP-violation effects in this scattering owing to the complex coupling constants. These exchange processes are active at some early epoch of the Universe, and they are out of equilibrium. In this case, at the relevant epoch, ordinary observer *Olga* should detect that

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º It should be specified that T_R, T'_R mean the temperatures at the moment when the energy density of relativistic products of the inflaton decay started to dominate over the energy density of the inflaton oscillation.

(i) matter slowly (in comparison to the Universe expansion rate) disappears from the thermal bath of O-world, (ii) particles and antiparticles disappear with different rates, and at the end of the day she observes that her world acquired a non-zero baryon number even if initially it was baryon symmetric.

On the other hand, her mirror colleague *Maxim* would see that (i) matter creation takes place in M-world, (ii) particles and antiparticles appear with different rates. Therefore, he also would observe that a non-zero baryon number is induced in his world.

One would naively expect that in this case the baryon asymmetries in the O- and M-sectors should be literally equal, given that the CP-violating factors are the same for both sectors. However, we show that in reality, the BA in the M sector, since it is colder, can be about an order of magnitude bigger than in the O sector, as far as washing out effects are taken into account. Indeed, these effects should be more efficient for the hotter O-sector while they can be negligible for the colder M sector, and this could provide reasonable differences between the two worlds if the exchange process is not too far from equilibrium. The possible marriage between dark matter and the leptobaryogenesis mechanism is certainly an attractive feature of our scheme.

The reactions relevant for the O-sector are the $\Delta L = 1$ one $l\phi \to \bar{l}'\bar{\phi}'$, and the $\Delta L = 2$ ones like $l\phi \to \bar{l}\bar{\phi}$, $ll \to \phi\phi$ etc. Their total rates are correspondingly

$$\Gamma_1 = \frac{Q_1 n_{\text{eq}}}{8\pi M^2} \; ; \quad Q_1 = \text{Tr}(D^\dagger D) = \text{Tr}[(y'^\dagger y')^* g^{-1}(y^\dagger y)g^{-1}]\,,$$

$$\Gamma_2 = \frac{3Q_1 n_{\text{eq}}}{4\pi M^2} \; ; \quad Q_2 = \text{Tr}(A^\dagger A) = \text{Tr}[(y^\dagger y)^* g^{-1}(y^\dagger y)g^{-1}]\,, \qquad (38)$$

where $n_{\text{eq}} \simeq (1.2/\pi^2)T^3$ is an equilibrium density per bosonic degree of freedom, and the sum is taken over all isospin and flavor indices of initial and final states. It is essential that these processes stay out of equilibrium, which means that their rates should not much exceed the Hubble parameter $H = 1.66 \, g_*^{1/2} T^2 / M_{Pl}$ for temperatures $T \leq T_R$, where g_* is the effective number of particle degrees of freedom, namely $g_* \simeq 100$ in the SM. In other words, the dimensionless parameters

$$k_1 = \left(\frac{\Gamma_1}{H}\right)_{T=T_R} \simeq 1.5 \times 10^{-3} \frac{Q_1 T_R M_{Pl}}{g_*^{1/2} M^2}\,,$$

$$k_2 = \left(\frac{\Gamma_2}{H}\right)_{T=T_R} \simeq 9 \times 10^{-3} \frac{Q_2 T_R M_{Pl}}{g_*^{1/2} M^2} \qquad (39)$$

should not be much larger than 1.

Let us now turn to CP-violation. In $\Delta L = 1$ processes the CP-odd lepton number asymmetry emerges from the interference between the tree-level and one-loop diagrams of fig. 3. However, CP-violation also takes place in $\Delta L = 2$ processes (see fig. 4). This is a consequence of the very existence of the mirror sector, namely, it comes from the contribution of the mirror particles to the one-loop diagrams of fig. 4. The direct calculation gives:[p]

$$\sigma(l\phi \to \bar{l}\bar{\phi}) - \sigma(\bar{l}\bar{\phi} \to l\phi) = \Delta\sigma \, ;$$
$$\sigma(l\phi \to \bar{l}'\bar{\phi}') - \sigma(\bar{l}\bar{\phi} \to l'\phi') = (-\Delta\sigma - \Delta\sigma')/2 \, ;$$
$$\sigma(l\phi \to l'\phi') - \sigma(\bar{l}\bar{\phi} \to \bar{l}'\bar{\phi}') = (-\Delta\sigma + \Delta\sigma')/2 \, ;$$
$$\Delta\sigma = \frac{3J\,S}{32\pi^2 M^4} \, , \qquad \Delta\sigma' = \frac{3J'\,S}{32\pi^2 M^4} \, , \tag{40}$$

where S is the c.m. energy square,

$$J = \operatorname{Im} \operatorname{Tr}[(y^\dagger y)^* g^{-1}(y'^\dagger y')g^{-2}(y^\dagger y)g^{-1}] \tag{41}$$

is the CP-violation parameter and J' is obtained from J by exchanging y with y'. The contributions yielding asymmetries $\mp\Delta\sigma'$ respectively for $l\phi \to \bar{l}'\bar{\phi}'$ and $l\phi \to l'\phi'$ channels emerge from the diagrams with $l'\phi'$ inside the loops, not shown in fig. 3.

Of course, this is in agreement with the CPT theorem that requires that the total cross sections for particle and anti-particle scatterings are equal to each other: $\sigma(l\phi \to X) = \sigma(\bar{l}\bar{\phi} \to X)$. Indeed, taking into account that $\sigma(l\phi \to l\phi) = \sigma(\bar{l}\bar{\phi} \to \bar{l}\bar{\phi})$ by CPT, we see that CP asymmetries in the $\Delta L = 1$ and $\Delta L = 2$ processes should be related by

$$\sigma(l\phi \to \bar{l}\bar{\phi}) - \sigma(\bar{l}\bar{\phi} \to l\phi) = \Delta\sigma \, ,$$
$$\sigma(l\phi \to X') - \sigma(\bar{l}\bar{\phi} \to X') = -\Delta\sigma \, , \tag{42}$$

where X' are the mirror sector final states, either $\bar{l}'\bar{\phi}'$ or $l'\phi'$. That is, the $\Delta L = 1$ and $\Delta L = 2$ reactions have CP asymmetries with equal intensities but opposite signs.

But, as L varies in each case by a different amount, a net lepton number decrease is produced, or better, a net increase of $B - L \propto \Delta\sigma$ (recall that

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[p] It is interesting to note that the tree-level amplitude for the dominant channel $l\phi \to \bar{l}'\bar{\phi}'$ goes as $1/M$ and the radiative corrections as $1/M^3$. For the channel $l\phi \to l'\phi'$ instead, both tree-level and one-loop amplitudes go as $1/M^2$. As a result, the cross section CP asymmetries are comparable for both $l\phi \to \bar{l}'\bar{\phi}'$ and $l\phi \to l'\phi'$ channels.
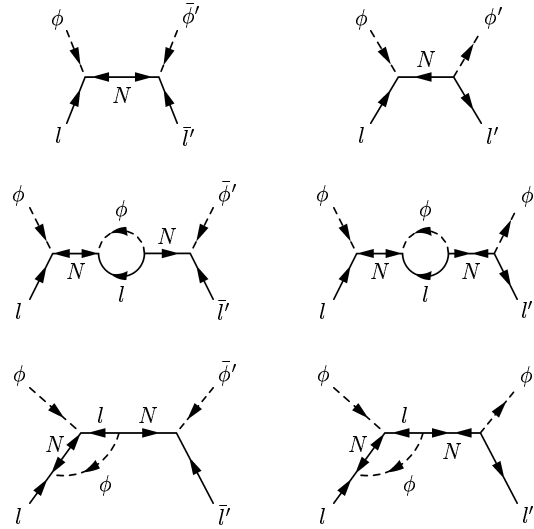
Figure 3.   Tree-level and one-loop diagrams contributing to the CP-asymmetries in $l\phi \to \bar{l}'\bar{\phi}'$ (left column) and $l\phi \to l'\phi'$ (right column).

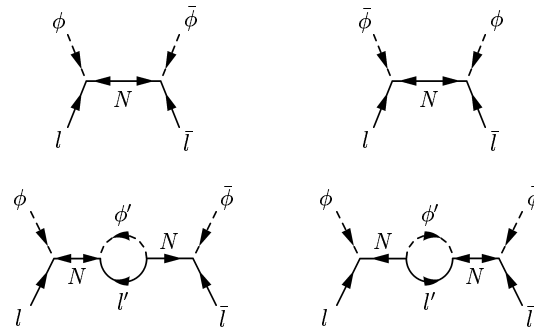
Figure 4.   Tree-level and one-loop diagrams contributing to the CP-asymmetry of $l\phi \to \bar{l}\bar{\phi}$. The external leg labels identify the initial and final state particles.

the lepton number L is violated by the sphaleron processes, while $B - L$ is changed solely by the above processes).

So long as we assume that the mirror sector is cooler and thus depleted of particles, the only relevant reactions are the ones with ordinary particles in the initial state. Hence, the evolution of the $B - L$ number density is determined by the CP asymmetries shown in Eqs. (40) and obeys the

equation

$$\frac{dn_{B-L}}{dt} + 3Hn_{B-L} + \Gamma n_{B-L} = \frac{3}{4}\Delta\sigma\, n_{\text{eq}}^2 = 1.8 \times 10^{-3}\,\frac{T^8}{M^4}\,, \qquad (43)$$

where $\Gamma = \Gamma_1 + \Gamma_2$ is the total rate of the $\Delta L = 1$ and $\Delta L = 2$ reactions, and for the CP asymmetric cross section $\Delta\sigma$ we take the thermal average c.m. energy square $S \simeq 17\,T^2$.

It is instructive to first solve this equation in the limit $k_{1,2} \ll 1$, when the out-of-equilibrium conditions are strongly satisfied and thus the term Γn_{B-L} can be neglected. Integrating this equation we obtain for the final $B - L$ asymmetry of the Universe, $Y_{BL} = n_{B-L}/s$, where $s = (2\pi^2/45)g_*T^3$ is the entropy density, the expression[q]

$$Y_{BL}^{(0)} \approx 2 \times 10^{-3}\,\frac{J\,M_{Pl}T_R^3}{g_*^{3/2}M^4}\,. \qquad (44)$$

It is interesting to note that $3/5$ of this value is accumulated at temperatures $T > T_R$ and it corresponds to the amount of $B - L$ produced when the inflaton field started to decay and the particle thermal bath was produced (Recall that the maximal temperature at the reheating period is usually larger than T_R.) In this epoch the Universe was still dominated by the inflaton oscillations and therefore it expanded as $a \propto t^{2/3}$ while the entropy of the Universe was growing as $t^{5/4}$. The other $2/5$ of (44) is produced at $T < T_R$, in the radiation dominated era when the Universe expanded as $a \propto t^{1/2}$ with conserved entropy (neglecting the small entropy leaking from the O- to the M-sector).

This result (44) can be recast as

$$Y_{BL}^{(0)} \approx \frac{20Jk^2T_R}{g_*^{1/2}Q^2M_{Pl}} \approx 10^{-10}\,\frac{Jk^2}{Q^2}\left(\frac{T_R}{10^9\text{ GeV}}\right)\,, \qquad (45)$$

where $Q^2 = Q_1^2 + Q_2^2$, $k = k_1 + k_2$ and we have taken again $g_* \approx 100$. This shows that for Yukawa constants spread e.g. in the range $0.1 - 1$, one can achieve $B - L = \mathcal{O}(10^{-10})$ for a reheating temperature as low as $T_R \sim 10^9$ GeV. Interestingly, this coincidence with the upper bound from the thermal gravitino production, $T_R < 4 \times 10^9$ GeV or so [38], indicates that our scenario could also work in the context of supersymmetric theories.

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[q] Observe that the magnitude of the produced $B-L$ strongly depends on the temperature; namely, larger $B - L$ should be produced in the patches where the plasma is hotter. In the cosmological context, this would lead to a situation where, apart from the adiabatic density/temperature perturbations, there also emerge correlated isocurvature fluctuations with variable B and L which could be tested with the future data on the CMB anisotropies and large scale structure.

Let us solve now Eq. (43) exactly, without assuming $\Gamma \ll H$. In this case we obtain [37]

$$Y_{BL} = D(k) \cdot Y_{BL}^{(0)} , \qquad (46)$$

where the depletion factor $D(k)$ is given by

$$D(k) = \frac{3}{5} e^{-k} F(k) + \frac{2}{5} G(k) \qquad (47)$$

and

$$F(k) = \frac{1}{4k^4} \left[(2k-1)^3 + 6k - 5 + 6e^{-2k} \right] ,$$

$$G(k) = \frac{3}{k^3} \left[2 - (k^2 + 2k + 2)e^{-k} \right] . \qquad (48)$$

These two terms in $D(k)$ correspond to the integration of (43) respectively in the epochs before and after reheating ($T > T_R$ and $T < T_R$). Obviously, for $k \ll 1$ we have $D(k) = 1$ and thus we recover the result as in (44) or (45): $= (B-L)_0$. However, for large k the depletion can be reasonable, e.g. for $k = 1, 2$ we have respectively $D(k) = 0.34, 0.1$.

Now, let us discuss how the mechanism considered above also produces the baryon prime asymmetry in the mirror sector. The amount of this asymmetry will depend on the CP-violation parameter $J' = \text{Im Tr}[(y^\dagger y)g^{-2}(y'^\dagger y')g^{-1}(y'^\dagger y')^* g^{-1}]$ that replaces J in $\Delta\sigma'$ of Eqs. 40). The mirror P parity under the exchange $\phi \to \phi'^\dagger$, $l \to \bar{l}'$, etc., implies that the Yukawa couplings are essentially the same in both sectors, $y' = y^*$. Therefore, in this case also the CP-violation parameters are the same, $J' = -J$.[r] Therefore, one naively expects that $n'_{B-L} = n_{B-L}$ and the mirror baryon density should be equal to the ordinary one, $\Omega'_b = \Omega_b$.

However, now we show that if the $\Delta L = 1$ and $\Delta L = 2$ processes are not very far from equilibrium, i.e. $k \sim 1$, the mirror baryon density should be bigger than the ordinary one. Indeed, the evolution of the mirror $B - L$ number density, n'_{B-L}, obeys the equation

$$\frac{dn'_{B-L}}{dt} + 3Hn'_{B-L} + \Gamma'n'_{\text{B-L}} = \frac{3}{4}\Delta\sigma' n_{\text{eq}}^2 , \qquad (49)$$

where now $\Gamma' = (Q_1 + 6Q_2)n'_{\text{eq}}/8\pi M^2$ is the total reaction rate of the $\Delta L' = 1$ and $\Delta L' = 2$ processes in the mirror sector, and $n'_{\text{eq}} = (1.2/\pi^2)T'^3 = x^3 n_{\text{eq}}$ is

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[r] It is interesting to remark that this mechanism needs the left-right parity P rather than the direct doubling one D. One can easily see that the latter requires $y' = y$, and so the CP-violating parameters J and J' are both vanishing.

the equilibrium number density per degree of freedom in the mirror sector. Therefore $k' = \Gamma'/H = x^3 k$, and for the mirror sector we have $Y'_{BL} = D(kx^3) \cdot Y^{(0)}_{BL}$. Hence, if $kx^3 \ll 1$, the depletion can be irrelevant, $D(kx^3) \approx 1$.

Now taking into account that in both sectors the $B - L$ densities are reprocessed into the baryon number densities by the same sphaleron processes, we have $B = a(B - L)$ and $B' = a(B - L)'$, with coefficients a equal for both sectors. Therefore, we see that the cosmological densities of the ordinary and mirror baryons should be related as

$$\beta = \frac{\Omega'_b}{\Omega_b} \approx \frac{1}{D(k)} \ . \tag{50}$$

If $k \ll 1$, depletion factors in both sectors are $D \approx D' \approx 1$ and thus we have that the mirror and ordinary baryons have the same densities, $\Omega'_{\rm b} \approx \Omega_{\rm b}$. In this case mirror baryons are not enough to explain all dark matter and one has to invoke also some other kind of dark matter, presumably cold dark matter.

However, if $k \sim 1$, then we would have $\Omega'_{\rm b} > \Omega_{\rm b}$, and thus all dark matter of the Universe could be in the form of mirror baryons. Namely, for $k = 1.5$ we would have from Eq. (50) that $\Omega'_b/\Omega_b \approx 5$, which is about the best fit relation between the ordinary and dark matter densities.

On the other hand, Eq. (45) shows that $k \sim 1$ is indeed preferred for explaining the observed magnitude of the baryon asymmetry. For $k \ll 1$ the result could be too small, since $(B - L)_0 \propto k^2$ fastly goes to zero.

One could question whether the two sectors would not equilibrate their temperatures if $k \sim 1$. As far as the mirror sector includes the gauge couplings which are the same as the standard ones, the mirror particles should be thermalized at a temperature T'. Once $k_1 \leq 1$, T' will remain smaller than the parallel temperature T of the ordinary system, and so the presence of the out-of-equilibrium hidden sector does not much affect the Big Bang nucleosynthesis epoch.

Indeed, if the two sectors had different temperatures at reheating, then they evolve independently during the expansion of the Universe and approach the nucleosynthesis era with different temperatures. For $k_1 \leq 1$, the energy density transferred to the mirror sector will be crudely $\rho' \approx (8k_1/g_*)\rho$ [20], where $g_*(\approx 100)$ is attained at the leptogenesis epoch. Thus, translating this to the BBN limits, this corresponds to a contribution equivalent to an effective number of extra light neutrinos $\Delta N_\nu \approx k/14$.

The following remark is in order. The mirror matter could be dark matter even if $k \ll 1$, when $Y'_{BL} = Y_{BL}$, if one assumes that the M-parity is

spontaneously broken so that the mirror nucleon masses are about 5 times larger than the ordinary ones. As we discussed in the previous section, this situation would emerge if $v'/v \sim 100$, and $\Lambda'/\Lambda \sim 5$. Needless to say, that the considered mechanism is insensitive to the values of the weak scale of the mirror sector so long as the latter remains much smaller than the masses of the heavy singlet neutrinos.

10. Mirror baryons as dark matter

We have shown that mirror baryons could provide a significant contribution to the energy density of the Universe and thus they could constitute a relevant component of dark matter. An immediate question arises: how the mirror baryon dark matter (MBDM) behaves and what are the differences from the more familiar dark matter candidates such as cold dark matter (CDM), hot dark matter (HDM) etc. In this section we briefly address the possible observational consequences of such a cosmological scenario.

In the most general context, the present energy density contains a relativistic (radiation) component Ω_r, a non-relativistic (matter) component Ω_m and the vacuum energy density Ω_Λ (cosmological term). According to the inflationary paradigm the Universe should be almost flat, $\Omega_0 = \Omega_m + \Omega_r + \Omega_\Lambda \approx 1$, which agrees well with the recent results on the CMBR anisotropy and large scale power spectrum.

The Hubble parameter is known to be $H_0 = 100h$ km s^{-1} Mpc^{-1} with $h \approx 0.7$, and for redshifts of cosmological relevance, $1 + z = T/T_0 \gg 1$, it becomes

$$H(z) = H_0 \left[\Omega_r (1+z)^4 + \Omega_m (1+z)^3 + \Omega_\Lambda \right]^{1/2}. \tag{51}$$

In the context of our model, the relativistic fraction is represented by the ordinary and mirror photons and neutrinos, $\Omega_r h^2 = 4.2 \times 10^{-5}(1 + x^4)$, and the contribution of the mirror species is negligible in view of the BBN constraint $x < 0.6$. As for the non-relativistic component, it contains the O-baryon fraction Ω_b and the M-baryon fraction $\Omega_b' = \beta \Omega_b$, while the other types of dark matter, e.g. the CDM, could also be present. Therefore, in general, $\Omega_m = \Omega_b + \Omega_b' + \Omega_{\mathrm{cdm}}$.[s]

The important moments for structure formation are related to the matter-radiation equality (MRE) epoch and to the plasma recombination and

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[s] In the context of supersymmetry, the CDM component could exist in the form of the lightest supersymmetric particle (LSP). It is interesting to remark that the mass fractions of the ordinary and mirror LSP are related as $\Omega'_{\mathrm{LSP}} \simeq x\Omega_{\mathrm{LSP}}$. The contribution of the mirror neutrinos scales as $\Omega'_\nu = x^3 \Omega_\nu$ and thus it is also irrelevant.

matter-radiation decoupling (MRD) epochs.

The MRE occurs at the redshift

$$1 + z_{\rm eq} = \frac{\Omega_m}{\Omega_r} \approx 2.4 \cdot 10^4 \frac{\omega_m}{1 + x^4} \tag{52}$$

where we denote $\omega_m = \Omega_m h^2$. Therefore, for $x \ll 1$ it is not altered by the additional relativistic component of the M-sector.

The radiation decouples from matter after almost all of the electrons and protons recombine into neutral hydrogen and the free electron number density sharply diminishes, so that the photon-electron scattering rate drops below the Hubble expansion rate. In the ordinary Universe the MRD takes place in the matter domination period, at the temperature $T_{\rm dec} \simeq 0.26$ eV, which corresponds to the redshift $1 + z_{\rm dec} = T_{\rm dec}/T_0 \simeq 1100$.

The MRD temperature in the M-sector $T'_{\rm dec}$ can be calculated following the same lines as in the ordinary one [19]. Due to the fact that in either case the photon decoupling occurs when the exponential factor in the Saha equations becomes very small, we have $T'_{\rm dec} \simeq T_{\rm dec}$, up to small logarithmic corrections related to B' different from B. Hence

$$1 + z'_{\rm dec} \simeq x^{-1}(1 + z_{\rm dec}) \simeq 1100\, x^{-1} \tag{53}$$

so that the MRD in the M-sector occurs earlier than in the ordinary one. Moreover, for x less than $x_{\rm eq} = 0.045\omega_m^{-1} \simeq 0.3$, the mirror photons would decouple during the radiation dominated period (see Fig. 5).

Let us now discuss the cosmological evolution of the MBDM. The relevant length scale for the gravitational instabilities is characterized by the mirror Jeans scale $\lambda'_J \simeq v'_s(\pi/G\rho)^{1/2}$, where $\rho(z)$ is the matter density at a given redshift z and $v'_s(z)$ is the sound speed in the M-plasma. The latter contains more baryons and less photons than the ordinary one, $\rho'_b = \beta\rho_b$ and $\rho'_\gamma = x^4\rho_\gamma$. Let us consider for simplicity the case when dark matter of the Universe is entirely due to M-baryons, $\Omega_m \simeq \Omega'_b$. Then we have

$$v'_s(z) \simeq \frac{c}{\sqrt{3}} \left(1 + \frac{3\rho'_b}{4\rho'_\gamma}\right)^{-1/2} \approx \frac{c}{\sqrt{3}} \left[1 + \frac{3}{4}\left(1 + x^{-4}\right)\frac{1 + z_{\rm eq}}{1 + z}\right]^{-1/2}. \tag{54}$$

Hence, for redshifts of cosmological relevance, $z \sim z_{\rm eq}$, we have $v'_s \sim 2x^2c/3 \ll c/\sqrt{3}$, quite in contrast with the ordinary world, where $v_s \approx c/\sqrt{3}$ practically until the photon decoupling, $z = 1100$.

The M-baryon Jeans mass $M'_J = \frac{\pi}{6}\rho_m\lambda'^3_J$ reaches the maximal value at $z = z'_{\rm dec} \simeq 1100/x$, $M'_J(z'_{dec}) \simeq 2.4 \cdot 10^{16} \times x^6[1 + (x_{\rm eq}/x)]^{-3/2}\omega_m^{-2}\ M_\odot$. Notice, however, that M'_J becomes smaller than the Hubble horizon mass $M_H = \frac{\pi}{6}\rho H^{-3}$ starting from a redshift $z_c = 3750x^{-4}\omega_m$, which is about $z_{\rm eq}$
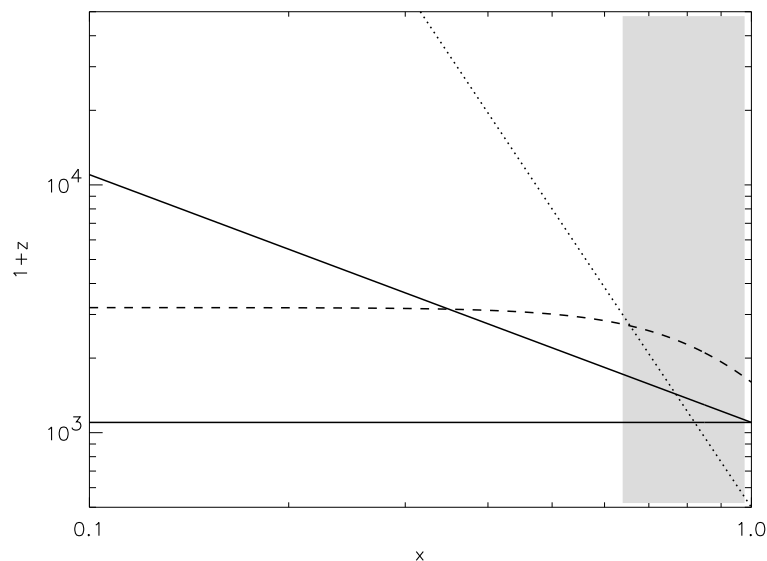
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Figure 5.   The M-photon decoupling redshift $1 + z'_{dec}$ as a function of x (thick solid line). The horizontal solid line marks the ordinary photon decoupling redshift $1 + z_{\text{dec}} = 1100$. We also show the matter-radiation equality redshift $1 + z_{\text{eq}}$ (dashes) and the mirror Jeans-horizon mass equality redshift $1 + z'_c$ (dots) for the case $\omega_m = 0.135$. The shaded area $x > 0.64$ is excluded by the BBN limits.

for $x = 0.64$, but it sharply increases for smaller values of x (see Fig. 5). So, the density perturbation scales which enter the horizon at $z \sim z_{\text{eq}}$ have mass larger than M'_J and thus undergo uninterrupted linear growth immediately after $t = t_{\text{eq}}$. The smaller scales for which $M'_J > M_H$ instead would first oscillate. Therefore, the large scale structure formation is not delayed even if the mirror MRD epoch did not occur yet, i.e. even if $x > x_{\text{eq}}$. The density fluctuations start to grow in the M-matter and the visible baryons are involved later, when after being recombined they fall into the potential wells of developed mirror structures.

Another important feature of the MBDM scenario is that the M-baryon density fluctuations should undergo strong collisional damping around the time of M-recombination. The photon diffusion from the overdense to underdense regions induces a dragging of charged particles and washes out the perturbations at scales smaller than the mirror Silk scale $\lambda'_S \simeq 3 \times f(x)\omega_m^{-3/4}$ Mpc, where $f(x) = x^{5/4}$ for $x > x_{\text{eq}}$, and $f(x) = (x/x_{\text{eq}})^{3/2}x_{\text{eq}}^{5/4}$ for $x < x_{\text{eq}}$.

Thus, the density perturbation scales which can undergo the linear growth after the MRE epoch are limited by the length λ'_S. This could help in avoiding the excess of small scales (of a few Mpc) in the power spectrum without tilting the spectral index. The smallest perturbations that survive the Silk

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damping will have the mass $M_S' \sim f^3(x)\omega_m^{-5/4} \times 10^{12}\ M_\odot$. Interestingly, for $x \sim x_{\text{eq}}$ we have $M_S' \sim 10^{11}\ M_\odot$, a typical galaxy mass. To some extent, the cutoff effect is analogous to the free streaming damping in the case of warm dark matter (WDM), but there are important differences. The point is that like usual baryons, the MBDM should show acoustic oscillations with an impact on the large scale power spectrum.

In addition, the MBDM oscillations transmitted via gravity to the ordinary baryons, could cause observable anomalies in the CMB angular power spectrum for l's larger than 200. This effect can be observed only if the M-baryon Jeans scale λ_J' is larger than the Silk scale of ordinary baryons, which sets a principal cutoff for CMB oscillations around $l \sim 1200$. As we have seen above, this would require enough large values of x, near the BBN upper bound $x \simeq 0.6$ or so.

If the dark matter is entirely built up of mirror baryons, large values of x are excluded by the observational data. For the sake of demonstration, on Fig. 6 we show the CMBR and LSS power spectra for different values of x. We see that for $x > 0.3$ the matter power spectrum shows a strong deviation from the experimental data. This is due to Silk damping effects which suppress the small scale power too early, already for $k/h \sim 0.2$. However, the values $x < 0.3$ are compatible with the observational data.

This has a simple explanation. Clearly, for small x the M-matter recombines before the MRE moment, and thus it should rather manifest itself as the CDM as far as the large scale structure is concerned. However, there still can be a crucial difference at smaller scales which already went non-linear, like galaxies. Then one can question whether the MBDM distribution in halos can be different from that of the CDM? Namely, simulations show that the CDM forms triaxial halos with a density profile too clumped towards the center, and overproduces small substructures within the halo. As for the MBDM, it constitutes a sort of collisional dark matter and thus potentially could avoid these problems, at least clearly the one related to an excess of small substructures.

As far as the MBDM constitutes a dissipative dark matter like the usual baryons, one would question how it can provide extended halos instead of being clumped into the galaxy as usual baryons are. However, one has to take into account the possibility that during the galaxy evolution the bulk of the M-baryons could fastly fragment into the stars. A difficult question to address here is related to the star formation in the M-sector, also taking into account that its temperature/density conditions and chemical contents are very different from the ordinary ones. In any case, fast star formation
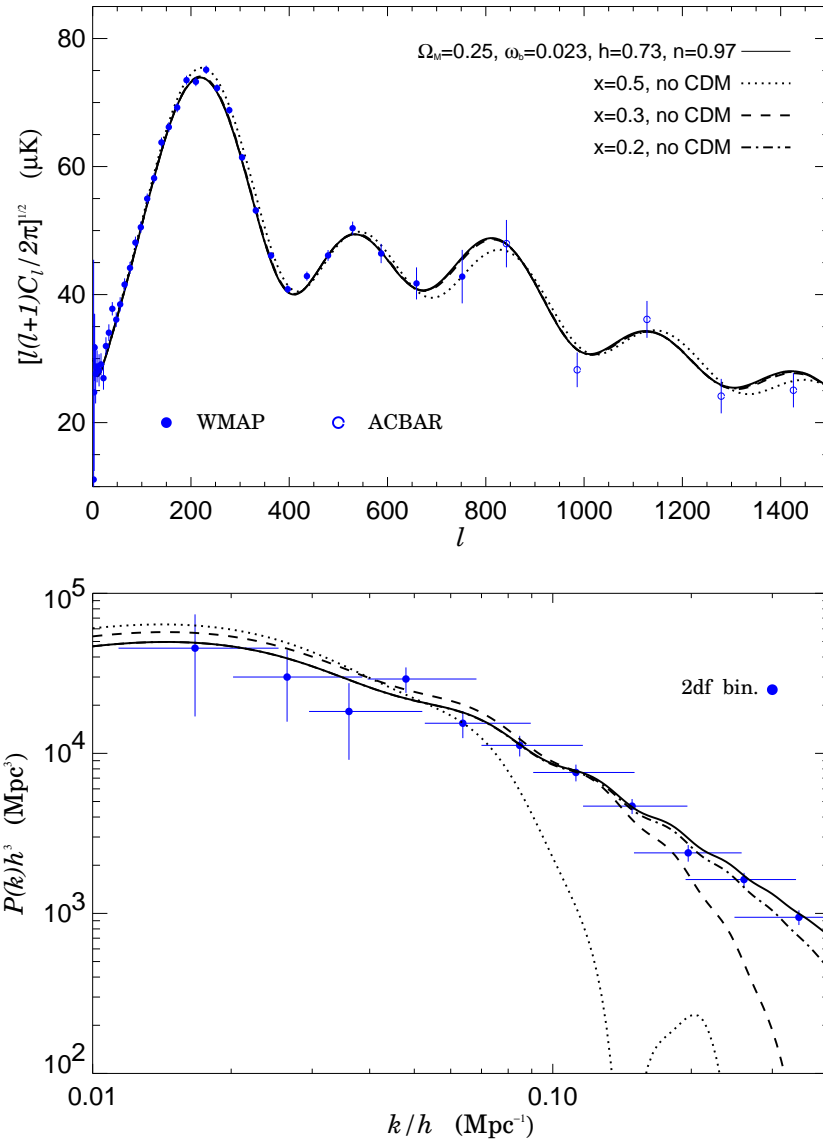
Figure 6.   The CMBR power spectrum (upper panel) and the large scale power spectrum (lower panel) for a "concordance" set of cosmological parameters. The solid curves correspond to the flat ΛCDM model, while dot, dash and dash-dot curves correspond to the situation when the CDM component is completely substituted by the MBDM for different values of x.

would remove the mirror gas and thus could avoid the M-baryons forming disk galaxies. The M-protogalaxy, which at a certain moment before disk

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formation essentially becomes a collisionless system of the mirror stars, could maintain a typical elliptical structure. In other words, we speculate on the possibility that the M-baryons form mainly elliptical galaxies.[t] Certainly, in this consideration also the galaxy merging process should be taken into account. As for the O-matter, within the dark M-matter halo it should typically show up as an observable elliptic or spiral galaxy, but some anomalous cases are also possible, like certain types of irregular galaxies or even dark galaxies dominantly made of M-baryons.

Another tempting issue is whether the M-matter itself could help in producing big central black holes, with masses up to $\sim 10^9 \ M_\odot$, which are thought to be the main engines of active galactic nuclei.

Another possibility that can also be considered is when dark matter in galaxies and clusters contains mixed CDM and MBDM components, $\Omega_d = \Omega'_b + \Omega_{cdm}$. e.g. one can exploit the case when mirror baryons constitute the same fraction of matter as the ordinary ones, $\Omega'_b = \Omega_b$, a situation which emerges naturally in the leptogenesis mechanism of Sect. 4.3 in the case of small k.

In this case the most interesting and falsifiable predictions are related to the large x regime. In Fig. 7 we show the results for the CMBR and LSS power spectra. We see that too large values of x are excluded by the CMBR anisotropies, but e.g. $x \leq 0.5$ can still be compatible with the data.

The detailed analysis of this effect will be given elsewhere [39]. In our opinion, in case of large x the effects on the CMBR and LSS can provide direct tests for the MBDM and can be falsified by the next observations with higher sensitivity.

In the galactic halo (provided that it is an elliptical mirror galaxy) the mirror stars should be observed as Machos in gravitational microlensing [14, 40]. Leaving aside the difficult question of the initial stellar mass function, one can remark that once one assumes the mirror stars could be very old and evolving faster than the ordinary ones, it is suggestive to think that most of the massive ones, with mass above the Chandrasekhar limit $M_{\rm Ch} \simeq 1.5 \ M_\odot$, have already ended up as supernovae, so that only the lighter ones remain as the microlensing objects. The recent data indicate the average mass of Machos around $M \simeq 0.5 \ M_\odot$, which is difficult to explain in terms of the brown dwarves with masses below the hydrogen ignition limit $M < 0.1 M_\odot$ or other baryonic objects [41]. Perhaps, this is the observational evidence of mirror matter?

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[t] For a comparison, in the ordinary world the numbers of spiral and elliptic galaxies are comparable. Remarkably, the latter contain old stars, very little dust and show low activity of star formation.
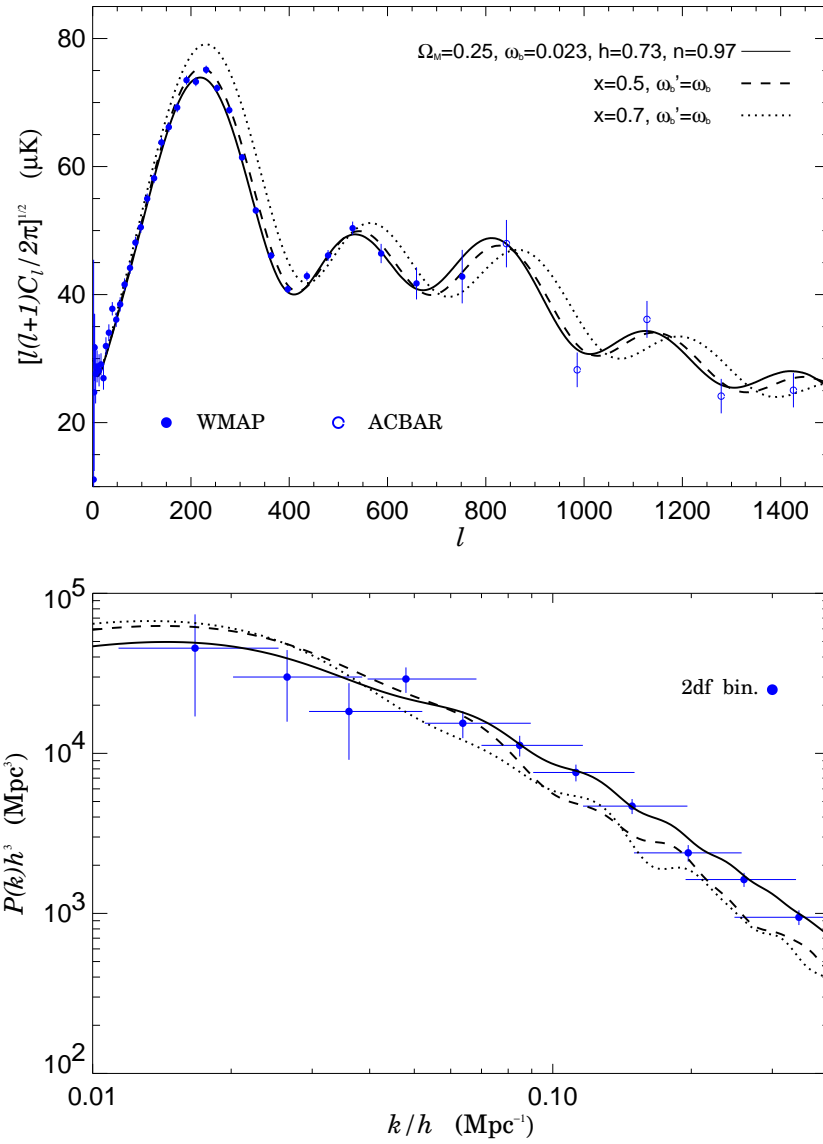
Figure 7.   The same as on Fig. 4, however for the mixed CDM+MBDM scenario for large values of x. The ordinary and mirror baryon densities are taken equal, $\Omega'_b = \Omega_B$, and the rest of matter density is attained to the SDM component.

It is also possible that in the galactic halo some fraction of mirror stars exists in the form of compact substructures like globular or open clusters. In this case, for significant statistics, one could observe interesting time and angular correlations between the microlensing events.

The explosions of mirror supernovae in our galaxy cannot be directly seen by an ordinary observer. However, they should be observed in terms of gravitational waves. In addition, if the M- and O-neutrinos are mixed [12,13], it can lead to an observable neutrino signal, and could also be accompanied by a weak gamma ray burst [42].

11. Conclusions and outlook

We have discussed cosmological implications of the parallel mirror world with the same microphysics as the ordinary one, but having smaller temperature, $T' < T$, with the limit on $x = T'/T < 0.6$ set by the BBN constraints. Therefore, the M-sector contains less relativistic matter (photons and neutrinos) than the O-sector, $\Omega'_r \ll \Omega_r$. On the other hand, in the context of certain baryogenesis scenarios, the condition $T' < T$ yields that the mirror sector should produce a larger baryon asymmetry than the observable one, $B' > B$. So, in the relativistic expansion epoch the cosmological energy density is dominated by the ordinary component, while the mirror one gives a negligible contribution. However, for the non-relativistic epoch the complementary situation can occur when the mirror baryon density is bigger than the ordinary one, $\Omega'_b > \Omega_b$. Hence, the MBDM can contribute as dark matter along with the CDM or even entirely constitute it.

Unfortunately, we cannot exchange information with the mirror physicists and combine our observations. (After all, as far as the two worlds have the same microphysics, life should be possible also in the mirror sector, and not only *Olga* but also *Maxim* could be a real person.) However, there can be many possibilities to disentangle the cosmological scenario of two parallel worlds with the future high precision data concerning the large scale structure, CMB anisotropy, structure of the galaxy halos, gravitational microlensing, oscillation of neutrinos or other neutral particles into their mirror partners, etc.

Let us conclude with two quotes of a renowned theorist. In 1986 Glashow found a contradiction between the estimates of the GUT scale induced kinetic mixing term (27) and the positronium limits $\varepsilon \leq 4 \times 10^{-7}$ and concluded that [8]: *"Since these are in evident conflict, the notion of a mirror universe with induced electromagnetic couplings of plausible (or otherwise detectable) magnitudes is eliminated. The unity of physics is again demonstrated when the old positronium workhorse can be recalled to exclude an otherwise tenable hypothesis"*.

The situation got another twist within one year, after the value $\varepsilon \sim 10^{-7}$ appeared to be interesting for tackling the mismatch problem of the or-

thopositronium lifetime. However, in 1987 Glashow has fixed that this value was in conflict with the BBN limit $\varepsilon < 3 \times 10^{-8}$ and concluded the following [9]: *"We see immediately that this limit on ϵ excludes mirror matter as an explanation of the positronium lifetime ... We also note that the expected range for ϵ $(10^{-3} - 10^{-8})$ is also clearly excluded. This suggests that the mirror universe, if it exists at all, couples only gravitationally to our own. If the temperature of the mirror universe is much lower than our own, then no nucleosynthesis limit can be placed on the mirror universe at all. Then it is also likely that the mirror universe would have a smaller baryon number as well, and hence would be virtually empty. This makes a hypothetical mirror universe undetectable at energies below the Planck energy. Such a mirror universe can have no influence on the Earth and therefore would be useless and therefore does not exist"*.

In this paper we objected to this statement. The mirror Universe, if it exists at all, would be useful and can have an influence if not directly on the Earth, but on the formation of galaxies ... and moreover, the very existence of matter, both of visible and dark components, can be a consequence of baryogenesis via entropy exchange between the two worlds. The fact that the temperature of the mirror Universe is much lower than the one in our own, does not imply that it would have a smaller baryon number as well and hence would be virtually empty, but it is likely rather the opposite, mirror matter could have larger baryon number and being more matter-rich, it can provide a plausible candidate for dark matter in the form of mirror baryons. Currently it seems to be the only concept which could naturally explain the coincidence between the visible and dark matter densities of the Universe. In this view, future experiments for direct detection of mirror matter are extremely interesting.

Perhaps mirror world is nothing but a reflection of our ignorance in the mirror of L-R equivalence. Perhaps Alice was wrong and there was no real Looking-Glass World beyond the mirror... However, its cosmology remains an interesting and non-trivial exercise for our imagination, which could help in understanding why our Universe looks as it is and how it could look under other circumstances...

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