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EVAPORATING GLOBAL CHARGES IN BRANEWORLD

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Dedicated to Stefan Pokorski on his 60th birthday

In braneworld models the global charges, such as baryon or lepton number, are not conserved. The global-charge non-conservation is a rather model-independent feature which arises due to quantum fluctuations of the brane worldvolume. These fluctuations create "baby branes" that can capture some global charges and carry them away into the bulk of higher-dimensional space. Such processes are exponentially suppressed at low-energies, but can be significant at high enough temperatures or energies. These effects can lead to a new, intrinsically high-dimensional mechanism of baryogenesis. Baryon asymmetry might be produced due either to evaporation into the baby branes, or creation of the baryon number excess in collisions of two Brane Universes.

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1. Introduction

The presently available theory of quantum gravity can be formulated in space-time with dimensionality greater than four. One possible scenario, of how our four-dimensional world emerges in this picture, is based on the assumption that all the Standard Model particles are localized on a 3-dimensional brane [1]. The absence of supersymmetry in the observable world could be related to a non-BPS nature of the brane [2]. Within the field theory context the simplest localization mechanism for fermions is due to the index theorem in the solitonic background [3]. As shown in [4], the localization of spin-1 gauge-fields in the field theory context is also possible but is more complicated. The mechanism requires the gauge group to be confining away from the brane [4]. This has an important model-independent consequence for the gauge-charge conservation on the brane, which is nothing but gauge flux conservation in the confining medium.

One of the motivations for the Brane Universe scenario is that it allows to lower the fundamental scale of quantum gravity all the way down to TeV or so, thus, providing a novel view on the Hierarchy problem [5,6]. The observed weakness of gravity at large distances is because gravitational fluxes spread into the N extra dimensions. The relation between the observed Planck scale, $M_{\rm P} \simeq 10^{19}$ GeV, and the so-called fundamental Planck scale M is then given by [5]:

$$M_{\rm P}^2 = M^{N+2} V_N \ , \qquad (1)$$

where $V_N \sim R^N$ with $N \geq 2$, is the volume of extra compactified spatial dimensions. The size of compactified radii in this picture can be in a sub-millimeter range without conflicting with any present astrophysical or laboratory constraints [7].

Below we argue that global charges are not conserved in the Brane Universe. The non-conservation of global charges is due to quantum fluctuations of the brane on which the Standard Model lives. These fluctuations can produce baby branes which can capture global charges and carry them away from the brane. At high enough temperatures or energies comparable with the brane tension the process of baby brane creation becomes significant. This leads to the global-charge transport from our brane. The corresponding process will look as non-conservation of global charges for a four-dimensional observer living on the brane[1]. These non-conservation mechanisms are significant in the cosmological context and can lead to new sources of baryogenesis. We discuss a possible cosmological scenario based on the Brane Inflation mechanism [11]. This scenario results in the baryon number access in our four-dimensional Brane Universe.

2. Non-conservation of global charges in the Brane Universe

In this section we discuss the mechanisms for non-conservation of global charges in the Brane Universe. The objective is to demonstrate that there are intrinsically high-dimensional phenomena which lead to baryon and lepton number non-conservation on our brane.

Let us start with the case when our $(3+1)$-dimensional Brane Universe is embedded in higher dimensional space-time. We will think of the brane

[1] These processes somewhat resemble the loss of quantum coherence in quantum gravity [8–10].

as being a dynamical object with the tension $\sigma \sim M^4$. This brane can fluctuate. The fluctuations are stronger at high temperature. In fact, there is a probability for the Brane Universe to wiggle strongly and create a baby brane. At high enough temperature the baby brane will be able to pull off the mother brane and propagate into the bulk of higher-dimensional space. This probability is non-vanishing if the temperature of the Brane Universe is nonzero (the same process could also be seen at high enough energies). The rate of this process is exponentially suppressed and can be estimated as [12]:

$$\frac{\Gamma}{\text{Volume}} \propto \exp\Big(-\frac{E_b}{T}\Big), \tag{2}$$

where E_b stands for the surface energy of the baby brane, which can be determined via its surface area A and tension σ, $E_b = A\sigma$. T stands for temperature of the Brane Universe. Thus, for high temperatures of order TeV or so, creation of baby branes should be an appreciable effect, while it should drop off rapidly as the brane cools down.

Once the baby brane is formed, it can capture some particles which happen to be nearby and carry them away from the mother brane. What will happen if the captured particles carry a net global charge, let us say baryon or lepton number? To clarify the issue, let us consider the case when the captured state is a gauge singlet combination of $u_{\rm R}, d_{\rm R}$ and $d_{\rm R}$ quarks (it might be any other combination of the Standard Model states which carries a non-zero baryon or lepton number but has *strictly zero* gauge charge). For a four-dimensional observer living on our brane this process will look as follows:

$$u_{\rm R} + d_{\rm R} + d_{\rm R} \to \text{NOTHING}, \tag{3}$$

where "NOTHING" stands for the baby brane which got separated from the mother brane and carries away the corresponding global charge. Since this object can gravitate, it will look for a four-dimensional observer as a piece of dark or hidden matter. Thus, the baby brane will carry its own baryon number $B_{\rm baby}$. The value of $B_{\rm baby}$ will exactly equal to the baryonic charge that is lost on the mother brane. If there are no bulk particles which can carry well defined baryon number, an observer on the mother brane will not be able to measure $B_{\rm baby}$. Thus, the process will look as disappearance of baryonic charge ΔB on the mother brane and appearance of the same charge $B_{\rm baby} = -\Delta B$ on the baby brane.

Let us try to understand this effect from the point of view of the effective four-dimensional field theory which is seen at distances larger than the size of extra dimensions. Let Ψ_M be a wave-function describing a state of some group of particles on our brane which carries a net baryonic charge. Likewise, we can define Ψ_B to be a wave-function of a set of particles on the baby brane. Both Ψ_M and Ψ_B are sharply localized functions in the bulk. The overlap between them is exponentially suppressed as the separation of the branes, r, in extra dimensions increases. Baryon number conservation in the theory is a result of the symmetry under the global phase rotations $\Psi_M \to e^{iQ_M\alpha}\Psi_M$, and $\Psi_B \to e^{iQ_B\beta}\Psi_B$, where Q_M and Q_B denote the charges for the corresponding wave-functions. When the branes are separated the wave-functions Ψ_M and Ψ_B get decoupled. As a result, there are two independent U(1) symmetries available: $\mathrm{U}(1)_M$ and $\mathrm{U}(1)_B$. However, when branes are close to each other, the functions Ψ_M and Ψ_B overlap. As a result, there is only one unbroken combination of $\mathrm{U}(1)_M$ and $\mathrm{U}(1)_B$, call it $\mathrm{U}(1)_{\mathrm{baryon}}$, which is defining baryon number of the whole system of the overlapping branes. In terms of the effective field theory language, the effective low-energy Lagrangian contains $\mathrm{U}(1)_M \otimes \mathrm{U}(1)_B$ violating term whose strength depends on the distance between the branes. This term drops exponentially fast as the separation between the branes increases:

$$(\bar{\Psi}_M)^{Q_B}\ (\Psi_B)^{Q_M}\ e^{-rM}. \tag{4}$$

What is important here is that the interaction should necessarily respect the $\mathrm{U}(1)_{\mathrm{baryon}}$ symmetry. As we mentioned above, for small r branes are interacting. Thus, there is only one baryon charge. This charge can be exchanged among the states of Ψ_M and Ψ_B. Suppose ΔQ denotes the amount of charge which is being transferred from the one set to another one. Once the branes are separated ($r \to \infty$) the overlap term disappears. Thus, there are two separately conserved charges corresponding to $\mathrm{U}(1)_M$ and $\mathrm{U}(1)_B$ respectively. However, only the charge Q_M will be seen in the mother brane and, thus, interpreted as the baryon number of our brane. Summarizing, the charge transport from Ψ_M to Ψ_B will look as disappearance of the ΔQ amount of the baryon charge on our brane and as appearance of exactly the same amount of the baryon charge on the baby brane.

Evidently, in each individual process ΔQ can take either sign and, if the system is in equilibrium, the net baryonic charge left on the brane will average to zero. However, it might be possible to generate a net baryon asymmetry on the mother brane if the system was *out of* equilibrium for some time during its evolution (it also requires C and CP violation [13], see discussions below). Below we address this issue and propose possible scenarios of how the baryon asymmetry could be generated in the Brane Universe. Before that let us discuss the fate of local charges in the Brane Universe. Seemingly, the same non-conservation process might be happening with the gauge charges, such as electric charge for instance. However, this cannot be true [4, 5]. Indeed, consider the case when the local charge is attached to the strongly fluctuating region of the mother brane which is about to be pulled off. The local charge, due to the corresponding flux conservation, would necessarily create a flux tube originating at the location of this charge and ending on the mother brane. At high enough energies, or temperatures likewise, the flux tube can break apart and the baby brane will eventually be liberated into the bulk of higher-dimensional space. However, the liberated baby brane will necessarily be neutral with respect to the local charge under consideration. Indeed, the process of breaking of the flux tube goes through creation of a charge-anti-charge pair in the tube. Once this pair is created, the anti-charge will get attracted by the original charge siting on the baby brane. Thus, the flux tube will break apart in such a way that the anti-charge from the pair will be attached to the baby brane and the charge of the pair will be attached to the mother brane. Hence, the final configuration of the liberated baby brane will be electrically neutral and the local charge will be conserved on the mother brane. Another way of saying this is to recall that all the Standard Model gauge interactions should be in a confining phase in the bulk space-time [4].

Summarizing, we conclude that the process of baby brane creation should lead to non-conservation of global charges (such as baryon or lepton number) in the Brane Universe. Moreover, this process will necessarily respect all the local charge conservation laws.

Other examples of such sharply localized objects can be bulk "glueballs" or "hadrons". These are the states that appear in the bulk due to the particular mechanism of localization of the gauge fields on the brane. As it was shown in [4], the field theory mechanism for localization of the massless gauge-fields on the brane implies that corresponding gauge group is in a confining phase in the bulk. Thus, a pair of test charges places in the bulk should be connected by a flux tube with the tension proportional to Λ^2, where Λ is a scale of the confining theory in the bulk. The inverse confinement scale, Λ^{-1}, sets the localization width for the observed gauge fields. For phenomenological reasons Λ should be greater than TeV.

Notice, that the gauge group in the bulk can be bigger than the Standard Model group. A photon, in this case, if being emitted into the bulk, becomes a gauge boson of the bigger confining theory. Thus, the photon can only escape the mother brane in the form of a heavy bound state, a sort of bulk "glueball". The similar consideration applies to fermion states. If the gauge group in the bulk were not confining, these fermions would have escaped the mother brane at energies bigger than the localization width. However, since the bulk is confining, such states can only escape within the corresponding "colorless" composite objects, bulk "Hadrons". Since the bulk "Hadrons" might carry off some net global charges, they can also lead to non-conservation of the global charges on the brane.

Below we study a field-theoretic model of non-conservation of global charges which is based on confining properties of the bulk gauge group.

As we mentioned before, the group is restored and confining outside of the brane [4]. For instance, the weak gauge group $\mathrm{SU}(2)_{\mathrm{L}}$ which is broken in our brane, should be restored (or be a part of a bigger unbroken group) and confining outside of the brane. The same applies to color $\mathrm{SU}(3)_c$ and hypercharge $\mathrm{U}(1)_{\mathrm{Y}}$ symmetries. $\mathrm{SU}(3)_c$ should either be a subgroup of bigger confining bulk gauge group, or be the same bulk gauge group with the confinement scale greater than TeV. Likewise, $\mathrm{U}(1)_{\mathrm{Y}}$ should be a part of a bigger group that is confining in the bulk. For simplicity of arguments we will be assuming that the gauge group within the brane is broken $\mathrm{SU}(2)_{\mathrm{L}}$ and outside of the brane it is confining $\mathrm{SU}(2)_{\mathrm{L}}$ (the generalization to the other groups and interactions is straightforward). Let us suppose that within the Brane Universe the confining phase which is realized outside of the mother brane is seen as a local false vacuum state. Then, in our four-dimensional world there is a finite probability to create a bubble (a sort of "hole") with a confining phase inside. If some "colorless" states with nonzero global charges are captured inside the bubble, they will be able to "leak" into the bulk. These effects are complimentary (but more model-dependent) to those discussed in the previous subsections. Let us study the bubble creation processes more carefully.

The probability to create a bubble per unit volume per unit time in our world with the confinement phase inside of the bubble is given by [14]:

$$\frac{P}{\mathrm{Volume}} \propto \exp\Big(- a \, \frac{\sigma^4}{(F(T) - \mathcal{E})^3}\Big) \ , \quad \text{when} \ \ \bar{\Lambda}^4 > F(T) > \mathcal{E}. \qquad (5)$$

Here, $F(T)$ denotes the free energy of the system as a function of temperature of the mother brane T, $\bar{\Lambda}^4$ denotes the depth of the scalar potential of the broken SU(2) theory, $\mathcal{E}$ is the difference between the energy densities of the confining and the Higgs phases, and a stands for some positive constant of order 10-100. As T is close to $\Lambda \sim$ few TeV this probability becomes significant. The theory inside of the bubble is in a confinement phase. Thus, bound states of particles which might form within the bubble are to be SU(2) singlets[2]. These singlet states will be able to propagate out of the Brane Universe. The most dramatic signature of this propagation is that they will be able to carry global quantum numbers off our Brane Universe. For instance, consider a single left-handed neutrino. This particle transforms in the fundamental dublet of $\mathrm{SU}(2)_{\mathrm{L}}$. Thus, it carries a "weak color" charge and cannot escape the brane. However, in accordance with 't Hooft's correspondence principle [15], the neutrino of the theory with a broken $\mathrm{SU}(2)_{\mathrm{L}}$

[2] If all the Standard Model interactions are considered, these states are supposed to be "color singlets" with respect to the whole Standard Model gauge group or w.r.t. the corresponding GUT, if the unification is assumed.

can be thought of as a "weak colorless" state, or as a bound state of confining $SU(2)_L$. Indeed, in the confinement picture, the left-handed neutrino can be presented as follows [15]:

$$\nu_L \text{ in Higgs phase } \langle = \rangle \ \bar{H}^i L_i \text{ in confinement phase} \ . \tag{6}$$

Here, H stands for the Standard Model Higgs dublet, $H_i^T = (\phi^+, \phi^0)$ and L stands for the left-handed dublet of a neutrino and electron, $L_i^T = (\nu_L, e_L)$. It is straightforward to see that the "weak colorless" bound state $\bar{H}^i L_i$ reduces to an ordinary left-handed neutrino once the Higgs field is given a non-zero vacuum expectation value (VEV). Indeed, in the unitary gauge $H^T = \frac{1}{\sqrt{2}}\Big(v + h, \ 0\Big)$, where v denotes the Higgs VEV and h stands for Higgs fluctuations about this VEV. Substituting this expression into the right hand side of Eq. (6) one finds, $\bar{H}^i L_i \to v \ \nu_L/\sqrt{2} + ...$ Thus, the r.h.s. of Eq. (6) can indeed be thought of as a "weak colorless" state of confining $SU(2)_L$; moreover, this state corresponds to the left-handed neutrino of the Standard Model.

Once the bubble is formed, the "weak colorless" state $\bar{H}^i L_i$ can appear in the confining phase inside of the bubble. This state, as we established above, carries leptonic charge. There is nothing that keeps this "weak colorless state" within the hot Brane Universe. Thus, it will be able to escape out into the higher-dimensional space. This process would seem as a leptonic charge non-conserving phenomenon to a four-dimensional observer living in the Brane Universe. The same applies to all the other standard model particles. Each of them can be thought of as "weak colorless" bound states [15]. Some of them are listed below:

$$\begin{aligned} e_L \text{ in Higgs phase } \langle = \rangle \ & \varepsilon_{ij} H^i L^j \text{ in confinement phase} \ ; \\ u_L \text{ in Higgs phase } \langle = \rangle \ & \bar{H}^i Q_i \text{ in confinement phase} \ ; \\ d_L \text{ in Higgs phase } \langle = \rangle \ & \varepsilon_{ij} H^i Q^j \text{ in confinement phase} \ ; \\ Z^0 \text{ in Higgs phase } \langle = \rangle \ & \bar{H} D_\mu H \text{ in confinement phase} \ . \end{aligned} \tag{7}$$

Here, Q denotes the left-handed up and down quark dublet. Some combinations of these states, such as (here we suppress all the Lorentz indexes and gamma matrices)

$$\varepsilon_{abc} \ \bar{H}^i Q_i^a d_R^b d_R^c, \quad \varepsilon_{abc} u_R^a d_R^b d_R^c, \tag{8}$$

will be created as "Standard Model colorless" excitations inside of those bubbles and, as a result, they will escape our brane at high enough temperatures or energies. Evidently, they will be able to carry the corresponding global charges, such as lepton or baryon number, away from the mother brane. This will make a four-dimensional observer think that the global quantum numbers are not conserved at high temperatures or energies in the Brane Universe. In the next section we discuss how these processes might lead to the baryon asymmetry in the Brane Universe.

3. Baryon asymmetry in the Brane Universe

In this section we argue that the baryon number non-conservation mechanisms discussed above might lead to a new approach to baryogenesis in the Brane Universe. We discuss two possible mechanisms. The first one is based on the fact that C and CP asymmetric branes can treat baryons and antibaryons differently. As a result, the rate to capture a baryon on a baby brane differs from that for an antibaryon. Thus, net baryon charge accumulation is possible if the system is out of equilibrium. The second scenario is based on production of the baryon number excess in a collision of two different Brane Universes after inflation. This scenario emerges naturally within the recently-proposed "Brane Inflation" framework [11].

As we discussed above, the baby branes and/or confining bubbles will carry some baryonic charge off our brane. The very same processes will be happening with antibaryons which will be taken away from the brane by the same mechanism. If the theory at hand does not distinguish between baryons and antibaryons, then the net charge carried away from our brane will average to zero. However, there is a possibility that the brane actually do distinguish between baryons and antibaryons if C and CP are broken. In particular, if the rate to capture a baryon on a baby brane differs from the corresponding rate for antibaryons, then the accumulation of the net baryonic charge on our brane will be possible in non-equilibrium processes [13].

Let us consider a toy model which demonstrates how this asymmetry can arise. Consider a scalar field χ which forms a four-dimensional "brane" embedded in five-dimensional space-time. Let us say the profile of this soliton is given by the familiar "kink" solution:

$$\chi = v \ \tanh(m x_5), \tag{9}$$

where m^{-1} defines the thickness of the brane and v stands for the VEV of the corresponding quantum field. Consider two five-dimensional fermions coupled to χ:

$$\mathcal{L}_{\rm int} = \chi \ (g_1 \bar{\psi}_1 \psi_1 + g_2 \bar{\psi}_2 \psi_2) + m_0 \ \psi_1^{\rm T} C^{(5)} \psi_2 \ + \ \text{other terms} \ , \tag{10}$$

where m_0 stands for some mass parameter and $C^{(5)}$ denotes the charge conjugation matrix in five-dimensional space-time, $C^{(5)} \equiv C \ \gamma_5$. This theory has the symmetry: $\psi_1 \to \exp(i\alpha)\psi_1$, $\psi_2 \to \exp(-i\alpha)\psi_2$. We identify this symmetry group with $\mathrm{U}(1)_{\mathrm{baryon}}$, thus ψ_1 and ψ_2 carry opposite baryonic charges. It is well known that each of these fermions give rise (in the massless limit) to a single chiral zero-modes localized on the brane:

$$\psi(x) \equiv \psi_1^0(x)\,\exp\Big(-\int_0^{x_5} g_1\chi(z)dz\Big),$$
$$\psi_c(x) \equiv \psi_2^0(x)\,\exp\Big(-\int_0^{x_5} g_2\chi(z)dz\Big). \tag{11}$$

From the point of view of the brane worldvolume field theory these chiral fermions can be identified with the worldvolume baryon ψ and antibaryon ψ_c (in Weyl notations)[3]. In the low-energy theory the "charge conjugation" symmetry $\psi \to \psi_c$ is broken since $g_1 \neq g_2$. This results in difference between the localization widths for ψ and ψ_c which are given by $\Delta \propto 1/g_1$ and $\bar{\Delta} \propto 1/g_2$ respectively. For instance, the width for (left handed) baryon can be made smaller than that for antibaryon ($g_1 > g_2$). Then, at energies $\Delta^{-1} < E < \bar{\Delta}^{-1}$ the antibaryon ψ_c can be "stripped off" the brane, while the baryon ψ would still be localized. This toy example explicitly shows how the brane can be "C-asymmetric". For generating net baryon charge, however, CP breaking is also required. Assuming that this is the case, (*i.e.* there are some explicitly CP-non-invariant terms in (10)), we expect that the probability for baryons to be captured by a baby brane is different than that for antibaryons (though, this process is more difficult to quantify). As a result, the baby branes will be able to remove from our world more antibaryons than baryons. Thus, the worldvolume observer will eventually see the net baryon asymmetry provided that "evaporation" into the baby branes is an out-of-equilibrium process. Such a out-of-equilibrium condition may emerge for instance from the reheating due to collisions of two Brane Universes.

Note that in this toy model there are bulk states which carry baryon number. They are Kaluza–Klein states of the original fermions ψ_1 and ψ_2. These states can mediate baryon number exchange between different branes. However, they are heavy, and the corresponding interactions are exponentially suppressed by the brane separation. Moreover, in realistic models due to the bulk confinement (which we have ignored in this toy example) these heavy states can only propagate within the bulk "colorless Hadrons".

Finally, we would like to discuss the issue of the over-closure of the Universe by baby branes in such a scenario. In order to generate the net

[3] Switching on small mass $m_0 \ll g_{1,2}\, v$ does not change the qualitative picture.

baryon asymmetry on our brane, not all the baby branes should return to it. If they stay in the bulk, they will look as a sort of dark matter with TeV mass. If we assume roughly one unit of baryon number captured per baby brane, their number density would be so large that they would over-close the Universe. However, there are several ways to avoid this problem. The most straightforward is to notice that the baby branes need not stay in the bulk, but rather can be "discharged" on some other distant brane (like ours, or even larger dimensionality). In such a case the energy density of baby branes will be converted into the distant brane tension and will be absorbed into the effective over-all cosmological term

$$\Lambda_{\rm eff} = \sum_i \sigma_i + \Lambda_{\rm bulk} V_N, \tag{12}$$

where $\Lambda_{\rm bulk}$ is the bulk cosmological constant and the summation is over all branes. The probability that the baby brane encounters a bigger brane and gets discharged there needs further quantification within more realistic models.

In this subsection we discuss the mechanism of baryogenesis which naturally arises within the Brane Inflation framework [11]. According to the general Brane Universe scenario, we live on a brane or a set of overlapping branes. The later possibility is supported by D-brane constructions in which the existence of a non-Abelian gauge group requires a number of parallel D-branes sitting on top of each other.

Before supersymmetry is broken branes are BPS states with zero net force between them. This is certainly true for two (or more) parallel D-branes, where the gravitational and dilaton attraction is exactly canceling with the repulsion mediated by Ramond–Ramond fields [16]. Similar examples can be constructed for field-theoretic branes, topological solitons [17]. However, in the real world supersymmetry must be broken and dilaton should be stabilized. Thus, we expect a non-zero net force between branes. The general expression for a potential between two such parallel branes embedded in $N > 2$ transverse dimensions at large distances ($r \gg M^{-1}$) takes the following form:

$$V(r) = M^4 \left(d + \frac{b_j {\rm e}^{-r m_j} - 1}{(M\ r)^{N-2}} \right). \tag{13}$$

The constant term d comes from the short-range brane-brane interaction. In fact, it accounts for interactions between particles localized on different branes, whose wave-functions only can overlap if branes intersect. The potential is normalized as $V(\infty) = 2\sigma$, σ being the brane tension. Yukawa potentials in (13) come from the exchange of heavy bulk modes with masses

m_j, and the power law interaction comes from the bulk gravitational attraction. If the D-brane picture is adopted, then m_j's should be understood as masses of dilaton and Ramond–Ramond fields. Regardless of what is the actual realization of branes, be it the D-brane picture or field theory soliton context, the potential in (13) describes adequately interactions between those objects. The model-dependent quantities are parameterized by coefficients d, b_j and m_j. These parameters determine the minimal separation $r_{\rm vac}$ at which the branes are stabilized in the lowest-energy state. If $r_{\rm vac} < M^{-1}$, the separation between the branes is smaller than the typical size at which the branes could fluctuate. Thus, the branes effectively sit on top of each other. As a result, the particles localized on these two branes are effectively shared by both of them. Below we will concentrate on the following alternative possibility. Let us assume that $r_{\rm vac} \gg M^{-1}$. In this case particles localized on two different branes have no overlap. Thus, they belong to either of branes, but are not shared among them. These two worlds can communicate to each other by exchanging bulk fields. If these interaction preserve global charges, B and L charges are conserved separately on each branes.

Let us see how this picture is affected by the dynamics of the brane inflation [11]. Once the branes are separated by a distance $r \gg r_{\rm vac}$, the nonzero potential energy between the branes gives rise to the four-dimensional effective cosmological constant that drives inflation [11]. This constant can be defined as follows:

$$\Lambda_{\rm eff} = V(r) + \Lambda_{\rm bulk} V_N, \tag{14}$$

where $\Lambda_{\rm bulk}$ is the bulk cosmological constant and V_N is the volume of the extra compactified space. Nearly zero value of the cosmological constant that is observed today implies that

$$\Lambda_{\rm vac} = V(r_{\rm vac}) + \Lambda_{\rm bulk} V_N \simeq 0. \tag{15}$$

Thus, according to Eqs. (13), (14), the four-dimensional vacuum cosmological constant will be nonzero for any $r \neq r_{\rm vac}$. This potential energy will drive inflation, the exponential growth of the three non-compact dimensions[2]. The next crucial thing is to note that for $r \gg r_{\rm vac}$ the potential (13) is a very flat function of r. As a result, the branes fall very slowly on each other. Thus, during this process the Universe is dominated by the potential energy which in fact triggers inflation in non-compact dimensions. We should also emphasis that the compact dimensions will not inflate since the effective Hubble size is never smaller than the size of the compact dimensions [11]. From the point of view of an effective four-dimensional theory this process

[2] The size of the extra dimensions will not be affected by this growth provided that the mass of the radius modulus is at least mm^{-1} [11].

is equivalent to slow rolling of a scalar field, an inflaton $\Phi = rM^2$. This field, according to (13) has a very flat potential. The quantity $\langle\Phi\rangle = r_{\rm vac}M^2$ is just the vacuum expectation value of the inflaton today.

The end of inflation is determined by the value of Φ which breaks either of the standard slow-roll conditions $V'M_{\rm P}/V < 1, V''M_{\rm P}^2/V < 1$ (see [11] for details). The epoch in which we are interested in starts right at this point of the evolution. We will argue below that after the branes collide and reheat each other, the net baryonic charge can be induced on our brane.

One possible scenario emerges when the branes get stabilized after the collision at some large distance $r_{\rm vac} \gg M^{-1}$. This is going to be the case if the branes repeal at short distances. For instance, this condition can be realized within the D-brane construction if dilaton becomes heavier than the corresponding Ramond–Ramond field $m_D \sim M \gg m_{\rm RR}$. As a consequence, when $r \ll m_{\rm RR}^{-1}$ the Ramond–Ramond repulsion takes over and branes get stabilized at $r_{\rm vac} \sim m_{\rm RR}^{-1}$.

Let us follow this scenario more closer. The potential energy of the Universe during inflation can be estimated as follows:

$$\Lambda_{\rm eff}(r \gg r_{\rm vac}) \sim M^4 \left(\frac{m_{\rm RR}}{M}\right)^{2-N}. \tag{16}$$

This amount of energy will transform into the energy of colliding branes after inflation. Let the wave-function of a set of particles localized on our brane be $\psi_{\rm our}(x_\mu)$, likewise, the wave function of a set of some different particles living on the other brane be $\psi_{\rm other}(y_\mu)$. There is a $\mathrm{U}(1)_{\rm our}$ baryon number symmetry on our brane, $\psi_{\rm our}(x_\mu) \to {\rm e}^{iQ_{\rm our}\alpha}\psi_{\rm our}(x_\mu)$. Likewise, there is a similar $\mathrm{U}(1)_{\rm other}$ symmetry on the other brane, $\psi_{\rm other}(x_\mu) \to {\rm e}^{iQ_{\rm other}\beta}\psi_{\rm other}(x_\mu)$. When branes are separated, these are two *different* symmetries. In the effective four-dimensional theory, this simply means that the interactions that break $\mathrm{U}(1)_{\rm our} \otimes \mathrm{U}(1)_{\rm other}$ are suppressed as follows:

$$(\psi^*_{\rm our})^{Q_{\rm other}}\ {\rm e}^{-rM}\ (\psi_{\rm other})^{Q_{\rm our}}. \tag{17}$$

However, once the branes come on top of each other, the suppression goes away. As a result, we are left with the only one conserved charge $Q = Q_{\rm other} + Q_{\rm our}$.

During inflation particles are inflated away on both branes and the expectation values of the operators $Q_{\rm other}$ and $Q_{\rm our}$ vanish. When the branes collide part of their energy is spent on creation of particles, baby branes and/or bubbles. Since the total charge Q is conserved, the net charge produced on the both branes should be zero. However, during the non-equilibrium collision process the branes overlap. Thus, $Q_{\rm other}$ and $Q_{\rm our}$ will *not* be separately conserved, and it might happen that in some reactions

$\Delta Q_{\rm our} = -\Delta Q_{\rm other} \neq 0$. Thus, the net global charges will be left on each branes. In addition to this effect, some charge will be carried away by the baby branes and/or the confining bubbles as discussed in the previous sections. We can briefly summarize the process described above as follows: The branes, while colliding, spend a very little time on top of each other. After that, they just "bounce back" and start to oscillate about the equilibrium point $r_{\rm vac}$. If C and CP symmetries are broken during the brane collisions, the couplings (17) allow "charges" to be "exchanged" among $\psi_{\rm our}$ and $\psi_{\rm other}$ during the short time moment of the collision. Thus, it might happen that one charge is produced in inflaton decays in excess and the other one in deficit (see the example below). Once the collision happened, these couplings switch-off almost instantly, and as a result, the values of nonzero charge asymmetries $\Delta Q_{\rm our} = -\Delta Q_{\rm other} \neq 0$ freeze-out. This, in particular, happens since the couplings (17) vanish almost instantly and the charges become separately conserved on two different branes.

The qualitative discussions given above can be made more precise by considering a simplified toy model. Consider two types of fermions, let us call them B_j and D_A. B_j's are localized on our brane and carry baryon number ($\mathrm{U}(1)_B$). D_A's, on the other hand, are localized on a distant brane and carry the corresponding global charge ($\mathrm{U}(1)_D$). Given the exponential suppression of the overlap of their wave-functions, a part of the effective four-dimensional Lagrangian for these fermions can be written as follows:

$$\begin{aligned}\mathcal{L}_{\rm int} \;=\; & c_{ij}(\Phi)\; B_i^c B_j \;+\; c_{AB}(\Phi)\; D_A^c D_B \;+\; \lambda_{iA}(\phi)\; \mathrm{e}^{-\frac{\Phi}{M}}\; D_A^c B_i \\ & + \bar{\lambda}_{iA}\; \mathrm{e}^{-\frac{\Phi}{M}}\; B_i^c D_A + c_{ijkm}(\Phi)\; B_i^c B_j B_k^c B_m \\ & + c_{ABCD}(\Phi)\; D_A^c D_B D_C^c D_D + \text{other interactions} \;+\; \text{H.c.}\,.\end{aligned} \tag{18}$$

Here, $\Phi \equiv M^2 r$ denotes a brane-separation modulus field, the inflaton, and c's and λ's are some polynomial functions of Φ in which C and CP violations are encoded. B_i^c and D_A^c stand for charge conjugated fields. Note that interactions of Φ with the fields on the same brane need not be exponentially suppressed. In some cases, these interactions arise after integrating out the bulk modes (*e.g.* open string modes stretched between two branes) which acquire masses due to the VEV of Φ and have direct couplings to the light modes on each brane. On the other hand, all the overlapping terms which break $\mathrm{U}(1)_B \otimes \mathrm{U}(1)_D$ symmetry explicitly must be exponentially suppressed (since, by the assumption, there are no light bulk modes with these charges).

Thus, when branes are well separated $\Phi \gg M$, the overlap terms are suppressed and the Lagrangian has two independent U(1)-symmetries. One of them acts on B's and can be regarded as baryon number symmetry in our brane. When branes come closer, however, the overlap terms do not vanish. As a result, we are left with one common fermion-number Abelian symmetry group $U(1)_F$. This last conserves the "total charge" of the branes $Q \equiv Q_B + Q_D$. Let us now turn to the particles which are being created in the inflaton decay. This decay, as we just mentioned, conserves the total charge Q. However, the individual charges, Q_B and Q_D are not conserved. Therefore, the rate for baryon number creation (*e.g.* in two-body decays) $\Phi \to B_i + D^c_A$, $\Phi \to B^{c*}_i + D^*_A$, and, likewise, the rate for antibaryon number creation, are different. Thus, although Q is conserved in the inflaton decays, individually Q_B and Q_D will not be conserved if both C and CP are broken.

Note that there might exist an additional source of physical CP violation due to "time interface" which can arise as a result of the time-dependent VEV of Φ and different dimensional operators present in the c and λ functions. These contributions are clearly very model-dependent and we will not attempt to quantify them here. An important outcome, however, is that in general, the rate to produce baryons in the inflaton decays differs from the same rate for antibaryons if C and CP are broken. Therefore, the nonzero value of $\Delta Q_B = -\Delta Q_D$ will be produced. When the branes bounce back after the collision, the inflaton VEV sharply increases. Thus, the $U(1)_B$ violating terms in (18) switch off and the baryon generation process stops before the system equilibrates. As a result, the accumulated net baryonic charge ΔQ_B freezes-out. Thus, our brane will be carrying the net baryon number [18] after the system comes to equilibrium.

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