Flux-Induced Baryon Asymmetry

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Abstract

I propose that the primordial baryon asymmetry of the universe was induced by the presence of a non-vanishing antisymmetric field background $H_{\mu\nu\rho}$ across the three space dimensions. This background creates a dilute $(B - L)$-number density in the universe cancelling the contribution from baryons and leptons. This situation naturally appears if the $U(1)_{B-L}$ symmetry is gauged and the corresponding gauge boson gets a Stuckelberg mass by combining with an antisymmetric field $B_{\mu\nu}$. All these ingredients are present in D-brane models of particle physics. None of the Sakharov conditions are required.

Dedicado a Leman
One of the most pressing cosmological puzzles is the observed baryon asymmetry in the universe. The traditional solution to this puzzle goes through the triplet of Sakharov conditions: a baryon asymmetry may be dynamically generated if the three ingredients 1) Baryon number violation, 2) C and CP-violation and 3) departure from thermal equilibrium take place in the history of the universe. A variety of concrete models which realize this general recipe have been proposed in the last 30 years. Although this recipe seems to work, one has the feeling that the way the baryon asymmetry appears in the history of the universe in this scheme depends very much on details of the models and is certainly not generic in particle physics models.

The philosophy underlying the idea of baryogenesis is that the primordial universe had the quantum numbers of the vacuum and hence it is natural to assume an exactly vanishing primordial baryon number and expect equal numbers of baryons and antibaryons. Assuming a small mismatch of order $10^{-10}$ for $n_B/n_\gamma$ as an initial condition is then totally unnatural.

The purpose of this note is to point out that the primordial universe could have the quantum numbers of the vacuum and still possess a primordial baryon-antibaryon asymmetry. The idea is that there could be a diluted distribution of baryon number density in the vacuum precisely cancelling the baryon number of baryons themselves. Specifically, I point out that under certain circumstances a constant non-vanishing antisymmetric field background $H_{\mu\nu\rho}$ across the three space dimensions may be such an extra source of baryon (or rather $B-L$ in the example discussed) number. In this scheme the overall primordial $B-L$ number vanishes. However, if a $H_{\mu\nu\rho}$ background is present, a net non-vanishing $B - L$ from baryons/leptons must be also present, due to $U(1)_{B-L}$ conservation. From this point of view having a $B-L$ asymmetry as an initial condition is something generic but still compatible with vanishing quantum numbers for the primordial universe. At lower temperatures electroweak instanton effects (violating the combination $B+L$) will force to have $n_B = -n_L$, but will be unable to erase the baryon and lepton asymmetries. Notice that the Sakharov conditions are not needed.

The essential idea is inspired by the generic phenomenon in D-brane string compacti-

1Perhaps an appropriate name for this could be ‘Baryonic Aether’. 
fications (for reviews and references see e.g. [1–4]) by which \( U(1) \) D-brane gauge bosons get Stuckelberg masses by combining with antisymmetric \( B_{\mu\nu} \) fields. In this case the \( U(1) \) symmetry survives as a global symmetry below the scale of the gauge boson mass. We here consider the case in which in addition there is a constant flux \( H_{\mu\nu\rho} \) along the three space dimensions. We find that in this case the latter gives rise to a density of \( U(1) \) charge proportional to the flux. A natural gauged \( U(1) \) symmetry to consider within the context of the SM is B-L, which is global symmetry of the SM. In the presence of 3 right-handed neutrinos this symmetry is anomaly-free and can thus be gauged without problems. This is in fact something generic e.g. in D-brane models of particle physics [1–4]. In such models the QCD gauge bosons come from three coincident parallel D-branes, giving rise to a gauge group \( U(3) = SU(3)_{QCD} \times U(1)_B \). The \( U(1)_B \) gauge symmetry corresponds to gauged baryon number. In the same way lepton number is the \( U(1) \) living on the worldvolume of a corresponding leptonic D-brane [5].

In fact a gauged \( U(1)_{B-L} \) appears also naturally in left-right symmetric extensions of the SM as well as \( SO(10) \) GUT models. However there is an important difference with our approach here. In the latter models the \( U(1)_{B-L} \) symmetry is assumed to be spontaneously broken by the standard Higgs mechanism. In here we are going to assume that the source for the mass of the \( U(1)_{B-L} \) gauge boson will be a Stuckelberg mass obtained by combining with an antisymmetric tensor \( B_{\mu\nu} \). Indeed, one interesting phenomenon observed while constructing SM-like intersecting D-brane models [5] is that often the gauge boson associated to the symmetry \( U(1)_{B-L} \) becomes massive by combining with an antisymmetric \( B_{\mu\nu} \) field. A coupling of type \( B \wedge F \) between the antisymmetric field \( B \) and the Abelian field strength \( F \) is the origin of this effect. As we review below, this term mixes both fields rendering the gauge boson massive by swallowing the \( B \) field. Note that in general this mechanism by itself does not give masses to neutrinos. We will assume that eventually some mechanism will give them either a \( (B-L) \)-preserving Dirac mass or else Majorana masses. As in the leptogenesis scenarios, for the observed values of neutrino masses, these Majorana masses are too small to erase the original B-L asymmetry.

It was known since a long time ago [6] that in string models there are pseudoanomalous
$U(1)$’s in which the triangle anomalies are cancelled by the exchange of some antisymmetric $B_{\mu\nu}$ field with appropriate couplings to the gauge bosons. This is the generalized Green-Schwarz mechanism [7] present in 4-dimensional string compactifications [6]. This works through the existence of two couplings, a $B \wedge F$ coupling as mentioned above and a coupling of type $\eta F \wedge F$, where $F$ is any gauge field strength in the theory and $\eta$ is the Poincare dual of $B_{\mu\nu}$. Under a $U(1)$ gauge transformation of parameter $\theta(x)$ it transforms like $\eta \to \eta + \theta(x)$. The combination of both terms renders the corresponding compactifications anomaly-free. In generic string models there may be a number of these ‘pseudo-anomalous $U(1)$s ’ which get anomaly-free making use of several antisymmetric fields $B$ [8]. As we said, the first of these terms, the $B \wedge F$ has the effect of rendering massive the corresponding pseudoanomalous $U(1)$. All pseudoanomalous $U(1)$’s in D=4 string compactifications become massive in this way. But the reverse is not true: an anomaly free $U(1)$ gauge boson (like B-L) may become massive if the appropriate $B \wedge F$ coupling is present. In ref. [5] it was shown that indeed those couplings are present in e.g., intersecting D6 brane models of particle physics. More recently it has also been found that in $N = 1$ SUSY models with MSSM-like spectra obtained from Type II rational CFT orientifolds, $U(1)_{B-L}$ gauge bosons do also get Stuckelberg masses in this fashion [9].

As I said the would be anomalous $U(1)_{B-L}$ gauge boson gets a mass of order the string scale $M_s$. On the other hand the corresponding $U(1)_{B-L}$ symmetry survives as a global rather than local symmetry in the low-energy lagrangian [10], [5,11]. This residual global $U(1)$ symmetries are rather generic in type II string D-brane models. It is not the case in traditional heterotic compactifications on CY with $SU(N)$ bundles. In the latter case there is a unique $B_{\mu\nu}$ field with indices in Minkowski which participates in the Green-Schwarz mechanism. A non-vanishing Fayet-Iliopoulos term is then induced [6] which forces some of the charged matter scalars carrying anomalous $U(1)$ charge to get vevs resulting in a breaking of the would-be global $U(1)$ symmetry. In the type II orientifold case such FI-terms vanish, as long as SUSY-breaking effects induce masses to scalars charged under the $U(1)$. Thus the global $U(1)$ symmetry remains perturbatively unbroken.

\[\text{\textsuperscript{2}}\text{In heterotic compactifications with } U(N) \text{ bundles instead of } SU(N) \text{ bundles [13] one gets again a}\]
Let us first recall how $U(1)$ gauge bosons gets massive before adding the H-flux. Consider the action of a gauge boson and an antisymmetric field $B_{\mu\nu}$ above. The relevant piece of the Lagrangian has the form

$$L_0 = -\frac{1}{12} H^{\mu\nu\rho} H_{\mu\nu\rho} - \frac{1}{4g^2} F^{\mu\nu} F_{\mu\nu} + \frac{cM}{4} \epsilon^{\mu\nu\rho\sigma} B_{\mu\nu} F_{\rho\sigma}$$  \quad (1)$$

where

$$H_{\mu\nu\rho} = \partial_\mu B_{\nu\rho} + \partial_\rho B_{\mu\nu} + \partial_\nu B_{\rho\mu}, \quad F_{\mu\nu} = \partial_\mu A_\nu - \partial_\nu A_\mu$$  \quad (2)$$

Here $M$ is a mass scale (of order the string scale $M_s$ in string models) and $c$ is a model dependent constant of order one. The last term in this expression is the $B \wedge F$ coupling mentioned above. Let us now review how the gauge boson gets mass by combining with the axion-like field $\eta$ (see e.g. [12]). We can rewrite this Lagrangian in terms of $H_{\mu\nu\rho}$ imposing the constraint $H = dB$ introducing a Lagrange Multiplier $\eta$:

$$L_0 = -\frac{1}{12} H^{\mu\nu\rho} H_{\mu\nu\rho} - \frac{1}{4g^2} F^{\mu\nu} F_{\mu\nu} - \frac{cM}{6} \epsilon^{\mu\nu\rho\sigma} H_{\mu\nu\rho} A_\sigma - \frac{cM}{6} \eta \epsilon^{\mu\nu\rho\sigma} \partial_\mu H_{\nu\rho\sigma}$$  \quad (3)$$

Now we can use the equations of motion for $H_{\mu\nu\rho}$ and find

$$H^{\mu\nu\rho} = -cM \epsilon^{\mu\nu\rho\sigma} (A_\sigma + \partial_\sigma \eta)$$  \quad (4)$$

Substituting back into eq.(3) one obtains

$$L_M = -\frac{1}{4g^2} F^{\mu\nu} F_{\mu\nu} - \frac{c^2 M^2}{2} (A_\sigma + \partial_\sigma \eta)^2$$  \quad (5)$$

This corresponds to the Stuckelberg Lagrangian of a massive vector boson of mass $cgM$.

Let us now proceed to the consideration of $H_{\mu\nu\rho}$ fluxes in a scheme with a gauged $U(1)$ as above. In principle such a vev would explicitly violate Lorentz invariance. However we will turn on fluxes only along the three space dimensions so that at the cosmological level there will be no contradiction with experimental facts. We will assume now a situation similar to that in type II D-brane models with several anomalous $U(1)$’s and axion-like fields.

In e.g. intersecting D6-brane models [1–5, 14] such 2-forms would be linear combinations of forms arising from wrapping the RR 5-form over 3-cycles in the CY. Then the flux considered would correspond to a RR flux $F_6$ with 3 legs on the 3-cycle and the other 3 on the space dimensions. Thus the $H_{\mu\nu\rho}$ flux here is not a NS flux.

Cosmology with a flux $H_{\mu\nu\rho}$ along the three space dimensions has been studied in the past, see e.g. [15] and references therein. However no connection with a baryon asymmetry was considered.
non-vanishing constant value of $H, h_{\mu\nu\rho}$. Then eq. (3) is modified to

$$L = L_0 - \frac{1}{12} h^{\mu\nu\rho} h_{\mu\nu\rho} - \frac{cM}{6} \epsilon^{\mu\nu\rho\sigma} h_{\mu\nu\rho} A_\sigma$$

and the final Lagrangian has the form

$$L = L_M - \frac{1}{12} h^{\mu\nu\rho} h_{\mu\nu\rho} - J_H A_\sigma - J_F A_\sigma$$

where

$$J_H^\sigma = \frac{cM}{6} \epsilon^{\mu\nu\rho\sigma} h_{\mu\nu\rho}$$

and we have added a term corresponding to the current $J_F^\sigma$ of the fermions coupling to the $U(1)$ gauge boson. In summary, we get a massive gauge boson but in addition the flux background $h_{\mu\nu\rho}$ acts like a current coupling to the massive $U(1)$ gauge boson. We will assume that the $H$-background is present in the universe only for the three spacelike components $x,y,z$ of the flux

$$h_{xyz} = H \epsilon_{xyz} \neq 0$$

so that actually the flux induces a $U(1)$-charge density. Although the $U(1)$ gauge boson is massive, we already pointed out that an unbroken global $U(1)$ symmetry persists. Thus, at the level of the low-energy effective Lagrangian (below the scale of the gauge boson mass) the effect of a vev for $H$ is to induce a non-vanishing global $U(1)$ charge density.

The above discussion applies to any gauged $U(1)$ symmetry whose gauge boson becomes massive a la Stuckelberg. This $U(1)$ may be anomaly free (like B-L) or anomalous, with the anomaly being cancelled by the Green-Schwarz mechanism. We want to apply these ideas to the case of the baryon number of the universe. Let us consider then for simplicity the case of a gauged $U(1)_{B-L}$ symmetry. In this situation we will have in the early universe a net vanishing B-L charge. Below the scale at which the $U(1)_{B-L}$ gauge boson gets a mass (i.e. the string scale in string models) (B-L)- number survives as a conserved global symmetry. Since this symmetry was gauged, in the primordial universe the overall B-L charge should vanish, very much like electric charge should vanish. Then the conservation of the residual global $U(1)_{B-L}$ current dictates that at the level of the effective field theory one has for the matter B-L density $n_{B-L}$

$$n_{B-L} + c MH = 0$$
Thus as long as we have $H \neq 0$ there will be a residual B-L number in the form of baryons/leptons and the ‘B-L aether’ induced by $H$ will compensate to get an overall vanishing (B-L) number. Note that both charges will be conserved separately. On the other hand at some point, at lower temperatures the electroweak instanton effects may be in thermal equilibrium giving rise to (B+L) violation in the standard way. The only effect of these will be to make $n_B = -n_L$ and the baryon-antibaryon asymmetry will persist.

One interesting question is what is the contribution of a non-vanishing $H$ to the present energy density in the universe. Is it sufficiently big to account for the observed cosmological constant? The answer is no, it contributes in a negligible way to the present energy density. Since $H$ and baryon densities are related by $|H| \simeq |n_B|/M$ one can easily make an estimate. As seen in eq.(6) one expects a contribution to the vacuum energy $V_H \propto H^2$ and hence one expects for the ratio of densities from $H$ and from baryons

$$\frac{\Omega_H}{\Omega_B} \propto \frac{H^2}{\rho_B} \propto \frac{\rho_B}{(m_P^2M^2)} \quad (11)$$

At present temperatures this ratio is extremely small for any reasonable value of the fundamental scale $M$ (one has $\Omega_H/\Omega_B \propto 10^{-80}$ for $M = 10^{16}$ GeV). Thus the contribution of the flux $H$ to the present vacuum energy density seems negligibly small. However, this contribution to the vacuum energy may have been much more important in the past. The reason for this is that [16] a background for $H$ in the Einstein’s equations behaves like ‘stiff matter’ (i.e., $p = \rho$) so that one has $\rho_H \propto 1/a(t)^6$, $a(t)$ being the scale factor. Compared to the baryon density one thus have

$$\frac{\rho_H}{\rho_B} \propto \frac{1}{a(t)^3} \quad (12)$$

Note that this behavior is consistent with equation (11) since it implies that $H^2$ scales with the scale factor like $n_B^2 \propto \frac{1}{a(t)^6}$, as expected. Thus the evolution equations are consistent with eq.(11) and the conservation of baryon number at any time.

The flux vev $H$ is in principle a free parameter of the underlying theory, very much like other fluxes considered recently in the context of string theory in order to stabilize the moduli [17]. An important difference is that these other fluxes go through the compactified extra dimensions whereas the flux here considered goes through the three space dimensions.
and has direct cosmological relevance. It would be interesting if we could figure out what is a natural value for the density $H$ since, given eq. (10), we could then compute the baryon asymmetry. It could well be that the density $H$ could be determined on anthropic grounds, certainly our existence very much depends on the amount of baryonic matter. On the other hand it would be interesting to have a specific model of string inflation in which the correct size for $H$ was dynamically determined. The $H$-background had to appear after inflation, otherwise it would have been totally diluted. If it was created say at the reheating temperature $T^*$, one can estimate the baryon asymmetry density then to be

$$\frac{n_B}{n_\gamma} \approx \frac{M H}{(T^*)^3}$$

(13)

If we insist in getting an asymmetry $n_B/n_\gamma \approx 10^{-10}$ one would need to have a flux

$$H \approx \frac{(T^*)^3}{M} 10^{-10}$$

(14)

If $M$ of order $10^{16}$ GeV (corresponding to a string scale of order the GUT scale) and a reheating temperature say of order $10^9$ GeV then the required flux at reheating is of order of a hadronic scale, $H \approx (300 \text{ MeV})^2$. If $M$ is of order the intermediate scale $M \approx 10^{11}$ GeV (as advocated in some string models), then the required flux is of order $H \approx (1 \text{ TeV})^2$. One can play around with different values for the reheating temperature and the string scale leading to the desired asymmetry.

If one could raise up the reheating scale close to a string scale of order $10^{16}$ GeV one could relate the asymmetry to the scale of SUSY-breaking. Although this sounds unlikely, let us explain it for the sake of the argument. Consider the context of type II orientifold string compactifications, which is a natural setting for the present mechanism. As we said, apart from this flux, in generic compactifications there are other antisymmetric fluxes (let me call them collectively $G$) which wrap cycles in the compact dimensions [17]. This is a crucial ingredient in recent efforts in order to understand the dynamical fixing of the string moduli. In addition to fixing the moduli, it has also been shown [18] that generically such fluxes $G$ do also break supersymmetry and give rise to SUSY-breaking soft terms of order $G/M_p$. In order to obtain soft terms of order $M_{sb} \propto 1 \text{ TeV}$ the fluxes must be diluted and be $G \propto M_{sb} M_p$. On the other hand at temperatures close to $T^* \approx M$ it is
natural to expect that both these fluxes and the one considered in this paper have similar densities, \( H \simeq G \simeq M_{sb}M_p \), since at that scale there is not much difference between compact and non-compact fluxes. If this is the case we would obtain for the asymmetry

\[
\frac{n_B}{n_\gamma} \simeq \frac{H}{M^2} \simeq \frac{M_{sb}M_p}{M^2}
\]  

(15)

For SUSY-breaking soft masses of \( M_{sb} \propto 1 \) TeV and \( M \) of order \( 10^{16} \) GeV (corresponding to a string scale of order the GUT scale) this ratio comes to be of order \( 10^{-10} \), of the order of the observed asymmetry. Note that if dark matter is constituted by SUSY neutralinos, both baryon asymmetry and dark matter would then be correlated since both would depend sensitively on the scale of SUSY-breaking. Let us however emphasize that this numerical exercise with such very high reheating temperature sounds unlikely within the present known models for reheating.

In this note we have emphasized the case of a gauged B-L symmetry because of its simplicity, it is anomaly-free and requires no Green-Schwarz mechanism. But is clear that the mechanism generalizes to the gauging of other possible global symmetries of the SM or any of its extensions, e.g. other linear combinations of B and L. The necessary ingredients are 1) a gauged \( U(1) \) symmetry (anomalous or not) whose gauge boson gets a mass term a la Stuckelberg and 2) a nonvanishing flux across the space dimensions for the corresponding antisymmetric field. Under those conditions the \( U(1) \) symmetry survives as a global symmetry and the flux induces an asymmetry of the corresponding charge. The only constraint is that the linear combination should be different from (B+L) because in this case electroweak instantons would erase any primordial asymmetry. One can also consider some other \( U(1) \)'s coming from some hidden sector of the theory, not coupling directly to the SM fields. In this case this mechanism could give rise to some density of hidden sector particles which could play the role of dark matter. If the relevant flux is of the same order of magnitude than that generating baryons and the masses of those dark matter objects is one order of magnitude larger than that of baryons, this could explain why dark and visible matter turn out to have not very different contributions to the energy of the universe.
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