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Baryogenesis in Supergravity Inflationary Models

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Realistic N=1 supergravity theories with a gravitino mass of order 1 Tev require a period of inflation to dilute the gravitino abundance. Moreover, if the gravitino is unstable the reheat temperature is bounded to be no greater than $O(10^8$ Gev). We show that such models may still have acceptable rates of baryosynthesis and discuss possible mechanisms.

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The standard big-bang cosmology imposes strong limits on the masses and lifetimes of unstable particles. These are particularly stringent for the gravitino whose couplings are only of gravitational strength leading to a very long lifetime $\tau = O(M^2/m^3_{3/2})$ which will decay after primordial nucleosynthesis. Here $M = 2.4\times10^{18}$ Gev and we often quote formulae with mass scale unit chosen to be M. Indeed these bounds rule out supersymmetric models with gravitino mass in the range $O(10 - 10^3)$ Gev, just that required in most realistic N=1 supersymmetric gauge field theories. In such models a period of inflation before baryogenesis is essential, and as a bonus may explain the flatness problem, the origin of small density fluctuations, the absence of monopoles etc.

Several groups have constructed supersymmetric inflationary models which [1,2,3] give adequate inflation to dilute the initial abundance of gravitinos to acceptable levels as well as giving the other attractive features of "the inflationary cosmology". However, one must beware of recreating a gravitino-like problem after inflation. This may happen by gravitino regeneration via particle collisions after inflation or by direct decay of the inflaton into the gravitino. (Similar problems may occur for production of any light, gravitino-like, weakly coupled state [4]). Excessive production through scattering processes may only be avoided if the reheat temperature, T_R is bounded above [5,6]. For unstable gravitinos the bound is[6]

$$T_R < 2.5 \times 10^8 \text{ Gev } \left(\frac{m_{3/2}}{100 \text{ Gev}} \right)^{-1} \qquad (1)$$

For stable gravitinos, in which the photino is heavier than the gravitino, the bound is less severe[6]

$$T_R < 1.3 \times 10^{12} \text{ Gev } \left(\frac{m_{3/2}}{100 \text{ Gev}} \right)^{-1} \qquad (2)$$

Most phenomenological models have been built using N=1 supergravity with gravitinos in the range 10-1000 Gev, in which it is usually assumed that photinos are lighter than gravitinos and bound (1) applies [For an alternative see ref [6]]. These models have difficulty in generating adequate baryon asymmetry. This is because the Higgs fields, whose baryon-number violating decays are responsible for producing the baryon-number asymmetry,can have substantial thermal production after inflation only if their masses are less than the reheat temperature T_R.

These Higgs boson fields are then so light that they lead to unacceptably fast proton decay rates in most SUSY Grand Unified Theories. However, in inflationary models baryogenesis may proceed via a novel mechanism in which the inflaton transfers its energy directly to these Higgs fields responsible for baryogenesis. These fields may have mass much greater than T_R - ie they need never be in thermal equilibrium and thus they are not Boltzmann suppressed in abundance. As a result they can efficiently produce a baryon asymmetry. In this letter we discuss in detail the implementation of this mechanism for inflationary models and the implications for T_R, the baryon asymmetry and proton decay.

We find acceptable levels of baryogenesis are possible in various models with reasonable proton decay rates. The first model uses a non-minimal Higgs sector and illustrates a novel mechanism in which an asymmetry in Massive-Higgs fields ultimately leads to a baryon asymmetry. The second model also uses an enlarged Higgs spectrum to generate acceptable fermion masses and to suppress dimension five proton decay amplitudes. In the latter the baryon-number and CP-violating interactions needed in baryogenesis involve right-handed neutrinos and thus do not contribute to proton decay. In both these models proton decay is expected with about a 10^{32} year lifetime, proceeding via scalar exchange. Finally, we briefly discuss baryogenesis in models in which the potential energy present during inflation is transferred to radiation and matter energy by anharmonic couplings between light and heavy fields and consider how this may lead to acceptable levels of baryogenesis.

In inflationary models the universe is reheated, after a period of exponential growth, by the out-of-equilibrium decay of a particle (or particles) denoted here by X. Then, following the analysis of Abbott, Farhi and Wise [7]

$$\frac{n_{B-\bar{B}}}{n_\gamma} \approx \frac{T_R}{m_X} \Delta B \qquad (3)$$

where m_X is the mass of X and ΔB is the baryon asymmetry produced per X decay.

To estimate $\frac{\Delta B}{m_X}$ it is first necessary to determine how potential energy is converted to matter and radiation energy after inflation. All supersymmetric models of inflation have so far been built using gauge singlet superfields which couple to the gauge non-singlet sector via gravitational interactions. The inflationary era is driven by a non-vanishing potential energy in the gauge singlet sector and, at the end of inflation, this potential energy is converted into energy in the gauge non-singlet sector reheating it to a temperature T_R. The detailed mechanism by which this happens depends on the coupling of the hidden to visible sectors. In this there is considerable freedom, but we will concentrate on models which inhibit the production of weakly coupled light scalars and avoid the "Polonyi problem" in which the energy density carried by these light states is unacceptably large [4].

The only known models which avoid the Polonyi problem are based on N=1 supergravity theories with minimal kinetic energy terms [3,8]. In these models the coupling between a hidden sector scalar field ϕ and the gauge non-singlet fields are given by two terms. First there is a term in the Lagrangian

$$e^{|\phi|^2/2} \quad [Q]_F \qquad (4)$$

where Q is the superpotential involving the gauge non-singlet fields, F denotes a $\Theta\Theta$ projection and we work in Planck units, of 2.4×10^{18} Gev. Secondly, there are purely scalar couplings coming from terms in the potential of the form

$$V = \left(\left| \frac{\partial P}{\partial \phi} + \phi^*(Q+P) \right|^2 + \left| \frac{\partial Q}{\partial \Sigma} + \Sigma^*(Q+P) \right|^2 - 3|Q+P|^2 \right) e^{\phi^*\phi + \Sigma^*\Sigma} \qquad (5)$$

where P is the superpotential of the hidden sector, Σ is any gauge non-singlet field, and all indices have been suppressed.

There are two ways these couplings may lead to the transfer of energy from the singlet sector. The first is through the direct decay of ϕ to gauge non-singlet fields lighter than ϕ. For illustration let us first consider the minimal form of Q in SU(5), employing the missing multiplet mechanism [9] to split the Higgs multiplets leaving heavy SU(3) triplets and light Higgs doublets. Then

$$Q = h_1\chi\chi H_1 + h_2\chi\psi H_2 + \alpha_1 H_1 \Sigma\Theta_2 + \alpha_2 H_2 \Sigma O_1 + M_O \theta_1\theta_2 + F(\Sigma)$$
$$+ h_3\theta_1\chi\chi \qquad (6)$$

where the representation assignments of the fields are given in Table 1 and $F(\Sigma)$ is a function chosen so that Σ acquires a vacuum expectation value (vev). Using eqs(4),(5) and (6) one may determine the possible decay modes of the inflaton to the gauge non-singlet sector. The dominant mode depends on the parameters in the theory.

For example, if H_1, H_2, θ_1 and θ_2 are heavier than ϕ (apart from the light Higgs doublet components H_1^D, H_2^D) then ϕ will decay via the couplings (cf eq(4) and (5))

$$e^{|\phi|^2/2}\,[h_1\chi\chi H_1^D + h_2\,\chi\psi H_2^D]_F \text{ and } [\phi^*(\frac{\partial P}{\partial\phi})^*]_A\,[h_1\chi\chi H_1^D + h_2\chi\psi H_2^D]_A \qquad (7)$$

(where A denotes the θ independent component) provided, of course they contain terms linear in ϕ. In general, this will be the case for usually ϕ acquires a very large vev of $O(M)$. If some of H_1, H_2, θ_1 and θ_2 are lighter than ϕ, there will be additional decay channels open giving decays to these states and the dominant one will depend on the relative magnitude of the Yukawa couplings and masses.

The second possible mechanism which can bleed energy from the ϕ sector was suggested by Ovrut and Steinhardt [2]. This mechanism involves an anharmonic coupling between ϕ and a more massive field Ψ. For example, an anharmonic term of the form $\phi^2\Psi$, in the potential energy, will drive oscillations in Ψ, even if Ψ is initially at rest, thus transferring energy from the light ϕ field to a heavy Ψ field which may decay to the gauge non-singlet sector by the couplings discussed above. Although Ovrut and Steinhardt considered an example in which Ψ was in the singlet inflation sector with direct couplings to ϕ in the superpotential, anharmonic terms may arise also through the couplings of eq(5) coupling ϕ to gauge non-singlet fields, such as Σ, which acquire a large vev. We will illustrate this possibility later.

Having identified the possible paths by which vacuum energy is released to reheat the universe after inflation we may now estimate $\frac{\Delta B}{M_X}$. For all of the mechanisms discussed reheat is achieved through out-of-equilibrium decays and so the first condition of Sakharov for a baryon asymmetry production is satisfied [10]. The second condition is that there should be baryon number violation in the reheating processes and this may occur either through direct production of baryon number violating Higgs or through their virtual exchange. Finally, CP violation is necessary in the baryon number violating process. This proves to be the most model-dependent aspect of the problem and we will illustrate several sources of adequate CP violation through explicit models.

Let us start our discussion with the minimal SU(5) model of eq (6) and direct inflaton decay. In this model it has been observed that H_1 or H_2 decay does not have the CP violation necessary to generate a baryon asymmetry. However, we will show that direct inflaton decay does produce an asymmetry. Consider two-body ϕ decay modes of Fig. 1. As we will show the gravitational couplings of eq (4) and (5) introduce the CP violation necessary to produce an asymmetry in θ_1 H_2, $\bar{\theta}_1\bar{H}_2$ production. If, in their decay, θ_1 and H_2 give, on average, states with baryon number B_{θ_1}, B_{H_2} then, if $B_{\theta_1} \neq - B_{H_2}$, the $\theta_1 H_2$, $\bar{\theta}_1\bar{H}_2$ asymmetry will be translated into a baryon number asymmetry.

For an asymmetry in $\theta_1 H_2$, $\bar{\theta}_1\bar{H}_2$ production it is necessary that the product of couplings in these graphs must be complex [11]. As as we have emphasised, there are two types of gravitationally induced couplings involving the inflaton and corresponding to eqs(4) and (5) respectively. Graphs involving the product of these couplings do have complex couplings and can generate a baryon asymmetry. This is clear from Fig 1(a) in which the right hand vertex corresponds to eq(4) while the left hand vertex corresponds to a term in eq(5) written, for convenience as a supersymmetric coupling $[\eta Q]_F$, where η is a spurion field with only a $\theta\theta$ component given by $F_{\eta_3} = [\,\phi^*(\frac{\partial P}{\partial\phi})^* + |\phi|^2 P^*\,]_A$ for the trilinear part of Q and $F_{\eta_2} = F_{\eta_3} - [P^*]_A$ for the bilinear part of Q. Fig 1a involves the couplings

$$|\alpha_{1,2}|^4 M_\Sigma \langle\Sigma\rangle^2 \,\langle \frac{\partial F_\eta}{\partial\phi} \rangle \quad \langle\phi\rangle \qquad (8)$$

and may be complex with phase α.

We will estimate the baryon asymmetry arising from Fig. 1 in the favoured limit of large M_Θ [9]. The states accessible in inflaton decay are $[H_1^P H_2^P]_F$, where H_1^P and H_2^P are the light eigenstates of mass $M_H \sim \frac{\alpha_1 \alpha_2 \langle\Sigma\rangle^2}{M_\Theta}$. In this case we estimate the baryon asymmetry generated by the graphs of Fig. 3 as

$$\Delta B \sim \frac{|\alpha_1|^2}{16\pi^2} \sin\alpha\ \delta\ \ \Delta b \qquad (9a)$$

where

$$\Delta b = 0\left(\frac{1}{10}\ \frac{h_2^2}{h_1^2}\ \frac{\Gamma_H}{M_H}\right) \text{ and } \delta = \frac{M_\phi}{M_\Sigma}\ \frac{M_H}{M_\Theta} \qquad (9b)$$

Δb is the net mean baryon number produced from the decay of H_1 and H_2 and Γ_H is their decay width. The couplings h_1 and h_2 generate top and bottom quark masses, so it is natural to take

$$\frac{h_1}{h_2} \sim \frac{m_t}{m_b}$$

giving $\Delta b = 0(10^{-4})$. In the SIC model with a reheat temperature $T_R \simeq 10^6$ Gev and $m_X \simeq 10^{10}$ Gev this gives an asymmetry $\Delta B = 0(10^{-10})$ provided $\delta \sin\alpha = 0(1)$ ie M_ϕ, M_H, M_Σ and M_Θ are nearly equal. While this could be arranged, such low values for M_Σ and M_Θ are very unnatural. This difficulty can be avoided by a minor modification of the theory which consists in adding a field A transforming as a singlet (or like the adjoint) representation of SU(5), and possessing the couplings $\alpha_3 H_1 H_2 A$ contributing in Fig (1b). In this case $\delta = \frac{M_\phi}{M_A}$ and for $\delta = 0(1)$ an adequate asymmetry is possible with M_Σ of the order of the GUT breaking scale, its natural value. Moreover in this version of SIC the reheat temperature is very low $T_R = \frac{\Delta^3}{M^2} \simeq 10^6$ GeV and comfortably within the bounds of eq (1).

Our example has shown that the new sources of CP violation involved in the gravitational couplings of eq (5) can lead to large baryogenesis.

However, there is a second difficulty, characteristic of all supersymmetric models, which must be avoided. Even though the Higgs triplets generating the baryon excess are heavier than the reheat temperature they cannot be arbitrarily heavy since their mass is necessarily lighter than the inflaton ϕ. In inflationary models the mass of ϕ is bounded if the density fluctuations, $\delta\rho/\rho$, are to be correctly given. For $\delta\rho/\rho < 10^{-4}$, $m_\phi < 10^{15}$GeV [12], and in simple inflationary models m_ϕ is much less than this bound. For example, in the SIC model [3], $m_\phi = \frac{\Delta^2}{M} \sim 10^{10} - 10^{11}$ GeV. This means that proton decay is expected to proceed rapidly through the graphs of Fig 2. Dimension 5 operators may be expected to dominate giving a proton decay rate

$$\tau_p = \frac{1}{x}\left(\frac{M_W}{\tilde{M}_W}\right)^2 \left(\frac{M_H}{10^{16}\text{GeV}}\right)^2 10^{32}\ \text{yr} \qquad (10)$$

Here we have assumed as is usual, $M_\Theta \gg \alpha_1 \langle\Sigma\rangle$ and we have included a factor x expressing our ignorance of the supersymmetric KM angles, $x \simeq 1$ in the minimal model. The factor $\left(\frac{M_W}{\tilde{M}_W}\right)$ is the ratio of the Wino Dirac mass to its Majorana mass component and is expected to be $\lesssim 10^2$ [13].
From eq(9) we see this value will not give adequate baryon asymmetry.

There are several ways this problem may be avoided. The estimate of the proton decay lifetime may be wrong because the relevant mixing angles are anomalously small giving a small value for x in eq (10). Alternatively the H_1H_2 mixing term may be small or zero. A third possibility is that the mass m_ϕ, which we took from the SIC model, may be larger allowing for heavier Higgs and thus reducing the proton decay rate. We will discuss these possibilities in turn, giving model estimates of the baryon asymmetry.

In ref [14] a model is presented in which the dimension five graphs of Fig. 2a contain a top quark and do not contribute to proton decay for kinematical reasons (ie. x=0). The form of eq (6) is maintained for the third generation, but further Higgs are added to generate the light fermion masses and mixing angles. In this model proton decay is mediated by the dimension six terms of Fig. 2(b) giving, for $m_{H_{1,2}} \approx 10^{10}$ Gev, a proton lifetime of $O(10^{31}$ yrs) with principal decay modes $p \to \bar{\nu}_\mu K^+$, $\mu^+ K^o$. The estimate of the baryon asymmetry produced by the direct inflaton decay graphs of Fig. 1(a), (b) still apply, but now there is no contraint on M_H following from proton decay. For $M_A = O(M_\phi)$ adequate baryogenesis is possible via the graph of Fig. 1(b).

In addition, as pointed out in ref [14], the new Higgs in this model can in introduce new sources of CP violation so that colour triplets produced from inflaton decay can themselves generate a baryon asymmetry in their decay via the graph of Fig. 3. However the asymmetry following from Fig. 3 is suppressed by the factor $(M_H/M_\Theta)^{3/2}$ and, for natural values of the parameters, will be smaller than the contributions from Fig. 1.
Let us consider now the second possibility for suppressing dimension five contributions by suppressing the H_1H_2 mixing term. We start with a generalization of the Georgi Jarlskog scheme [15] to allow for good predictions for all fermion masses in SU(5) and mixing angles, and to account for the recently measured lifetime of the B meson. The part of the superpotential which generates the $\Delta I_W = \frac{1}{2}$ masses is taken to be of the form [16].

$$G_1 = F^T \begin{bmatrix} 0 & AH & 0 \\ A'H & CL & 0 \\ 0 & 0 & BH \end{bmatrix} T + T^T \begin{bmatrix} 0 & DK & 0 \\ DK & EK' & FK'' \\ 0 & FK'' & GK \end{bmatrix} T \qquad (11)$$

A summary of the representation assignments of the fields is given in table 2. This form for the superpotential may be obtained by the imposition of the phase symmetries of Table 2. The dangerous dimension -five operators of Fig 2 are absent if there are no mass mixing terms of the form HK, HK', HK" etc and these are forbidden by the X symmetry of Table 2 which we identify as the (anomalous) Peccei-Quinn symmetry, which must be broken around 10^{10} Gev. Thus, in this model, proton decay does not proceed via the dimension -five graphs of Fig 2a.

The baryon number violation necessary for baryogenesis may be introduced in a way that does not lead proton decay, by coupling new massive states, S_i, most conveniently chosen to be matter-singlet superfields which can be identified with right-handed neutrinos à la SO(10). For this we note that the K,K" contain colour triplets which couple to matter in channels with different baryon number via superpotential terms allowed by the symmetries of table 2

$$G_2 = (aF_1S_2 + a'F_2S_1 + bF_3S_3)K + (cF_2S_3 + c'\,F_3S_2)K'' + (j_1S_1S_2 + j_2S_3S_3)\phi_1 \qquad (12)$$

This structure fits nicely into an SO(10) framework (Invariance under SO(10) would fix a=a' etc, and so we take them approximately equal). The supergraphs responsible for the baryon asymmetry are given in Fig 4 and correspond to heavy Higgs/Higgsino decay, where we assume these states have been produced via inflaton decay as discussed above.

Taking $a \approx a' \approx b \approx c \approx c'$ in eq(12); all Higgs masses to be the same order of magnitude, and, with Yukawa couplings of $O(10^{-2})$ leaves more than enough freedom in the Higgs sector to obtain $\Delta B = 10^{-5}$ to 10^{-6} which, in the S.I.C [3] model gives, via eq(1), $\frac{n_{B-\bar{B}}}{n_\gamma} \sim 10^{-10}$. Proton decay in this model may only proceed via the dimension -six Higgs scalar exchange graphs of Fig 2(b). Ref. [16] gives results for the dominant proton decay modes and bounds on the Higgs masses for various possible Higgs exchanges contributing to proton decay. The resulting lower limit on the (lightest) Higgs mass is consistent with the upper limits following from the need for adequate baryogenesis. For example in the SIC model $m_H < m_\phi \sim 10^{11}$ Gev. We see in this model there is an upper bound for the proton lifetime which is only one order of magnitude larger than the present experimental value.

Finally, let us discuss the third possibility that the mass of the field, X, whose decay reheats the universe, is much larger than in the SIC model. As discussed in ref [12], in general the inflaton mass cannot be heavier than $10^{-3}M_p$ if fluctuations $\delta\rho/\rho$ are to be $\leq O(10^{-4})$. It is just possible, if the inflaton saturates this bound, to have, in the minimal supersymmetric model, a long proton lifetime, (cf eq(10)) while generating an adequate baryon asymmetry via the graphs of Fig. 1 with the necessary CP violation coming from gravitational couplings of the inflaton as discussed above. Because the ϕ mass is so large, adequate baryogenesis could be obtained even with natural values for M_Σ and M_Θ; there is no need for the new A field. In order to obtain a heavy ϕ one needs to choose a more complicated inflaton potential than in the simple SIC model, but we see no reason, in principle, why this could not be done.

A more likely possibility is that the field X is <u>not</u> the inflaton field but is a much more massive field coupled via an anharmonic term to the inflaton. In the example of Ovrut and Steinhardt [2] discussed above X is the Ψ field with mass $O(M_P)$. Since it is so heavy its decay products may be very heavy too, allowing us to escape from the need for light, baryon number violating, Higgs and avoiding the problems associated with fast proton decay. However, we note that the CP violation necessary for baryogenesis is not present in the minimal model with couplings described by eq(6), because the vev of the Ψ field is very small and the graphs of Fig. 1 consequently give too small an asymmetry. Thus, in the Ovrut Steinhardt model, it is necessary to go to the non-minimal schemes such as discussed in our two examples given above. In both models adequate baryogenesis is possible but note that because M_X is $O(M_P)$ now there is no stringent upper bound on the Higgs mass and, consequently, no predicted upper bound on the proton lifetime.

To summarise: we have estimated the baryon excess which may be produced in inflationary models based on minimally coupled supergravity theories. Despite the fact these models necessarily have a low reheat temperature to avoid excessive thermal production of gravitinos, direct inflaton decay provides a mechanism which may produce adequate baryon asymmetry, provided the model has a source of large CP violation.

We found, even in the minimal supergravity model, that adequate CP violation can come from gravitationally-induced couplings. However, the need to avoid too fast proton decay requires that either the inflaton be anomalously heavy, or that the minimal Higgs structure be extended. In the latter case we find a lower bound for proton decay close to the current experimental limits. Another possibility has been suggested in which the inflaton transfers its energy via an anharmonic term to a massive scalar field. In this case baryogenesis would be possible in non-minimal models, but without a lower bound for proton decay.

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Figure Captions

Fig. 1. CP violating phase from gravitational couplings giving a baryon asymmetry in (a) the minimal supergravity model and (b) a model including a singlet or adjoint field A. In these figures the field R is given by Re $[\phi' < \phi^* >]$ where ϕ' is the propagating part of ϕ.

Fig. 2. Supergraphs leading to proton decay (a) via dimension five operators and (b) via dimension six operators. In fig (2a) λ denotes a gaugino.

Fig. 3. Supergraphs giving a baryon asymmetry in the model of ref [13].

Fig. 4. Baryogenesis induced by the decay of colour triplet Higgses into massive right handed neutrinos.

Table 1

multiplet	χ	ϕ	H_1	H_2	θ_1	θ_2	Σ	A
representation	10	$\bar{5}$	5	$\bar{5}$	50	$\overline{50}$	75	1 or 24

Table 2

multiplet	F_1	T_1	F_2	T_2	F_3	T_3	H	L	K	K'	K"	S_1	S_2	S_3	ϕ_1
representation	$\bar{5}$	10	$\bar{5}$	10	$\bar{5}$	10	$\bar{5}$	$\overline{45}$	5	5	5	1	1	1	1
V	-3	1	-3	1	-3	1	2	2	-2	-2	-2	5	5	5	-10
X	1	1	1	1	1	1	-2	-2	-2	-2	-2	1	1	1	-2
F	1/2	1/2	-1/2	-1/2	0	0	0	1	0	1	1/2	1/2	-1/2	0	0

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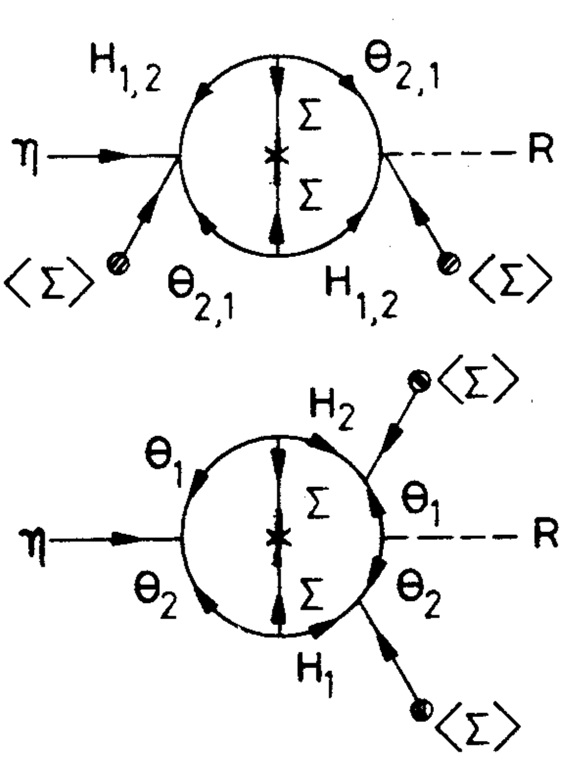

FIG. (1a)

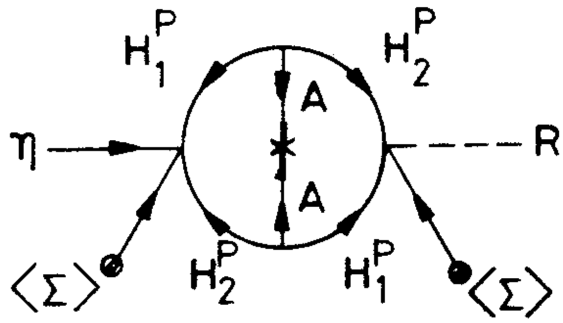

FIG. (1b)

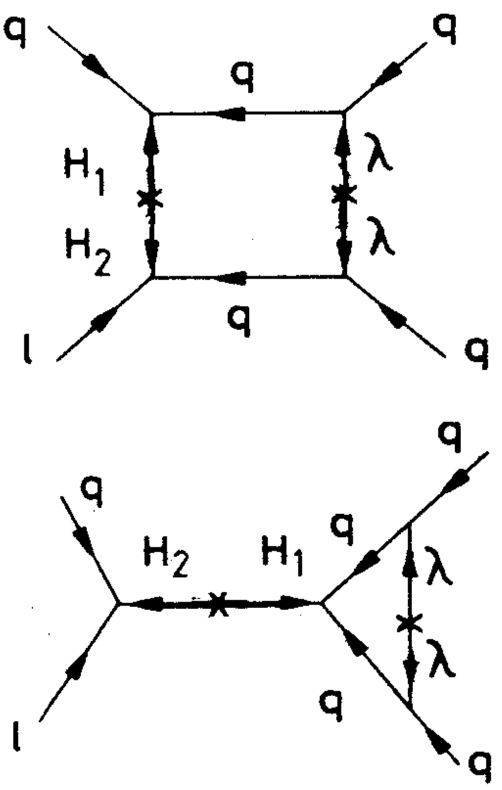

FIG. (2a)

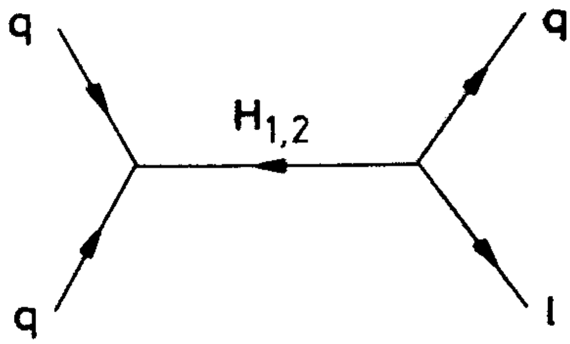

FIG (2b)

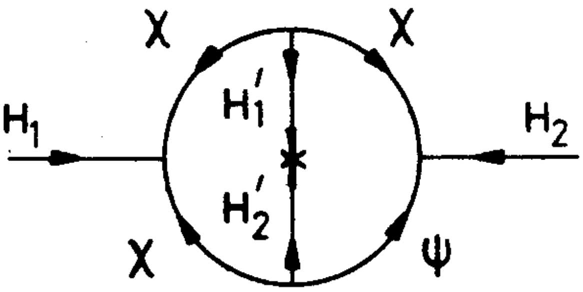
χ
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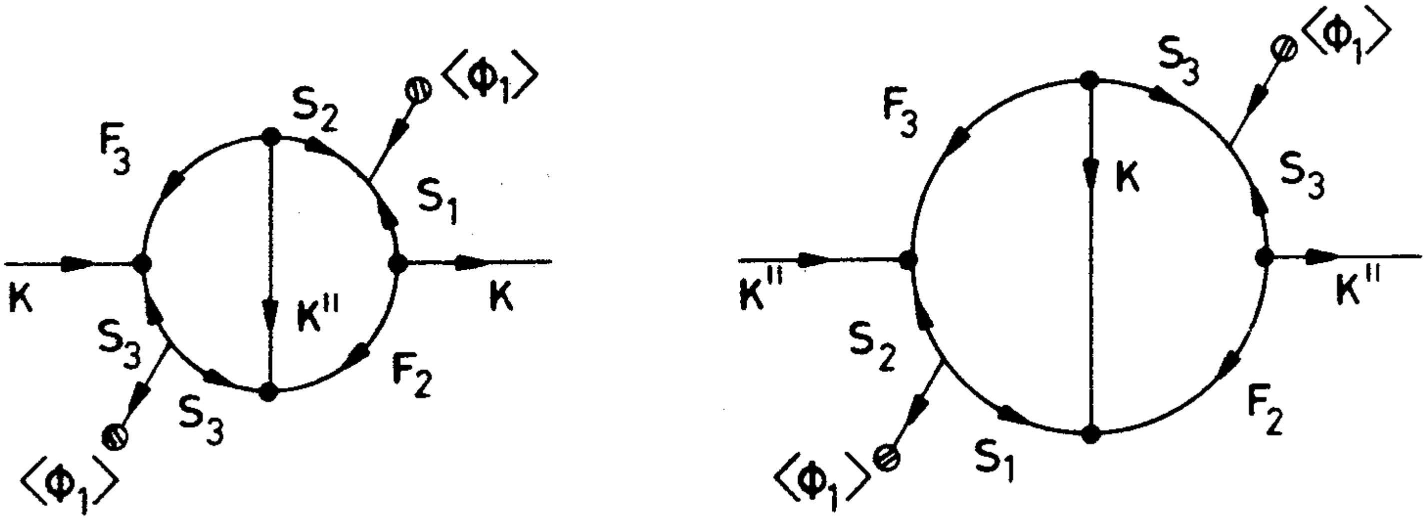
FIG. 3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K
K''
K
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K
S₃
K''
K''
S₂
F₂
⟨Φ₁⟩
S₁

FIG. 4