DEUTSCHES ELEKTRONEN-SYNCHROTRON DESY

DESY 87-026
March 1987

ADJUSTING THE COSMOLOGICAL CONSTANT DYNAMICALLY: COSMONS AND A NEW FORCE WEAKER THAN GRAVITY

by

R.D. Peccei, J. Solà, C. Wetterich

Deutsches Elektronen-Synchrotron DESY, Hamburg

ISSN 0418-9833

NOTKESTRASSE 85 · 2 HAMBURG 52

To be sure that your preprints are promptly included in the
HIGH ENERGY PHYSICS INDEX ,
send them to the following address (if possible by air mail) :

DESY
Bibliothek
Notkestrasse 85
2 Hamburg 52
Germany

DESY 87-026
(Revised)
May 1987

ISSN 0418-9833

Adjusting the cosmological constant dynamically: Cosmons and a new force weaker than gravity

R.D. Peccei, J. Solà[1] and C. Wetterich
Deutsches Elektronen-Synchrotron DESY
Hamburg, Fed. Rep. Germany

Abstract

We argue that the vanishing of the cosmological constant obtains as a result of the dynamics of a new field - the cosmon - which is the Goldstone boson of dilatation invariance, assumed to be broken spontaneously near the Planck scale. The presence of the cosmon, coupled to the fact that scale invariance is anomalous at the quantum level, drives the cosmological constant to zero, provided that the energy momentum trace is purely anomalous. Furthermore, these quantum anomalies give the cosmon a small mass, giving rise to an intermediate range force ($\lambda \leq 10^4 m$). We can estimate the effect of the cosmon force between macroscopic matter distributions. The dominant component of the force is attractive, couples to mass and should be weaker than gravity ($\alpha \sim 10^{-2} - 10^{-3}$). There is, however, also a repulsive baryon number dependent component of calculable strength ($\sim 1/20\alpha$) and an even smaller contribution proportional to Z-N.

[1]On leave from Departament de Física, Universitat Autònoma de Barcelona, Barcelona, Catalonia.

Einstein's equations admit Minkowski space as a solution only if the trace of the energy momentum tensor $T^{\mu\nu}$ in vacuum vanishes ($< T^{\mu}_{\mu} >$ is the cosmological constant). In the standard model of the strong and electroweak interactions, however, there exist various nonvanishing contributions to this trace. For instance, chiral symmetry breaking in QCD gives a contribution to $< T^{\mu}_{\mu} >$ proportional to $< m_u \bar{u}u + m_d \bar{d}d > \simeq -f_{\pi}^2 m_{\pi}^2 \simeq -(1.7 \times 10^{-4} GeV^4)$, while the gluon condensate gives a contribution almost two orders of magnitude greater $< \frac{\beta(g_s)}{2g_s} F_a^{\mu\nu} F^a_{\mu\nu} > \simeq -(1.4 \times 10^{-2} GeV^4)$ [1]. The contribution to $< T^{\mu}_{\mu} >$ of the Higgs sector, unless it is appropriately tuned to zero, is of order $10^8 GeV^4$. Why should these contributions, plus possibly others arising from sectors of the theory beyond the standard model, add up to zero? [F1] This is the long standing cosmological constant problem [2].

Since the problem arises within physics at scales much below the Planck mass, M_P, we think that its solution should also find a description within an effective "low energy" theory. Here we investigate the possibility that the cosmological constant problem is solved by a new interaction, whose range is sufficiently larger than the inverse QCD scale Λ_{QCD}, so that long range coherent effects lead to a dynamical adjustment of the cosmological constant to zero. In many ways the mechanism we suggest is similar to the invisible axion solution [3] to the strong CP problem. In the axion case the imposition of an extra symmetry [4], which is broken at a high scale, suffices to dynamically cause the vanishing of the parameter $\bar{\theta}$, responsible for the strong CP breaking. We shall argue here that something quite analogous happens for the cosmological constant.

A natural candidate for the quantum mediating this new interaction is that it be a scalar field, which is a singlet with respect to the low energy gauge symmetry $SU(3) \times SU(2) \times U(1)$. Such a singlet can couple to ordinary quarks, leptons and gauge fields only via effective nonrenormalizable (dimension 5 or higher) interactions. These are naturally suppressed by some mass scale. If we assume that this scale is of order M_P, one would understand why such a scalar was not discovered in ordinary particle physics experiments. However, this new interaction could well compete with gravity, if its range was large enough, giving rise to testable consequences. As we shall see, these latter phenomenological aspects are one of the interesting consequences of our proposal.

There is a different, more theoretical, motivation for considering interactions involving a scalar singlet, which is connected with dilatation symmetry. As is well known [5], the divergence of the dilatation current is linked to the trace of the energy momentum tensor. Thus the cosmological constant problem has a natural relation to the fate of dilatation symmetry. Our principal assumption is that, at the fundamental level, one has an action which is invariant under coordinate scalings: $x \rightarrow e^{-\alpha} x$, accompanied by corresponding scalings of the quantum fields, according to their dimensions: $\Phi \rightarrow e^{\alpha}\Phi; \Psi \rightarrow e^{3/2\alpha}\Psi$; etc. [F2]. Since we know that a mass scale, the Planck mass, appears connected with gravity, this dilatation symmetry must be realized in a Nambu Goldstone manner, being spontaneously broken at a scale M near M_P. At low energies, the dilatation symmetry of the theory is not manifest, except through the couplings of the Goldstone mode of dilatation symmetry to ordinary matter and gauge fields. It is these interactions which will serve to drive the cosmological constant to zero.

It is easy to write down a dilatation symmetric version of the action of the standard model. One introduces for this purpose the Goldstone mode, S, of dilatation symmetry, which translates under scale transformations: $S \rightarrow S + \alpha M$, where M is the scale of the spontaneous breakdown of dilatation symmetry (typically $M \sim M_P$). Then a dilatation symmetric action is obtained by multiplying any parameter with dimension $(mass)^D$ with a factor $e^{DS/M}$. In the standard model these parameters are only found in the Higgs sector. We must, therefore, replace the Higgs mass parameter μ^2 by $\mu^2 e^{2S/M}$ and replace a possible constant term κ in the Higgs potential by $\kappa e^{4S/M}$. In addition we must add a scale invariant kinetic energy for the S field:

$$\mathcal{L}_S = 1/2(\partial^\mu S)(\partial_\mu S)e^{2S/M} \tag{1}$$

At the classical level the standard model Lagrangian, augmented in the manner indicated above, possesses a conserved dilatation current. This is easily seen to be:

$$J^\mu = Me^{2S/M}\partial^\mu S + \sum_i D_i \chi_i \frac{\partial \mathcal{L}}{\partial(\partial_\mu \chi_i)} + x_\nu T^{\mu\nu} \tag{2}$$

where D_i is the dimension of the various fields χ_i entering in the theory. Because the divergence of the first two terms on the RHS of Eq. (2) vanishes in the vacuum, by Lorentz invariance, it follows that also $< T^\mu_\mu >$ vanishes. However, in most cases the vanishing of $< T^\mu_\mu >$ is trivial, resulting from S being driven to $-\infty$, so that all scale parameters in the theory vanish! For a realistic theory one needs $< T^\mu_\mu >= 0$ and a finite static value for S in the vacuum. Since the theory is scale invariant, there will be no classical static solution for the field S unless the parameters κ and μ^2 in the Higgs potential are correlated. For S to have a static solution $< \frac{\partial V}{\partial S} >$ must vanish. This will only happen if the Higgs potential takes the form

$$V(\Phi, S) = \lambda[\Phi^\dagger\Phi - \Phi_0^2 e^{2S/M}]^2 \tag{3}$$

in which case a static solution will exist for all values of S [F3]. Clearly, if (3) holds it is obvious that the trace of the classical energy momentum tensor vanishes in the vacuum

$$< T^\mu_\mu >= 4 < V >= M < \frac{\partial V}{\partial S} >= 0 \tag{4}$$

This also obtains for more general parameters for which $< \frac{\partial V}{\partial S} >$ does not vanish automatically. However, as we already mentioned, in that case S is driven to minus infinity and $< T^\mu_\mu >= 0$ only because the field Φ cannot obtain a vacuum expectation value.

For a discussion of the cosmological constant we should of course not restrict our discussion a priori to flat space. This is easily remedied by adding a dilatation invariant gravity piece to the Lagrangian of the theory

$$\mathcal{L} = -\sqrt{g}\frac{M_P^2}{16\pi}e^{2S/M}R \equiv -\sqrt{g}hM^2 e^{2S/M}R \tag{5}$$

The model admits then additional solutions with constant finite expectation values for Φ and S and constant curvature scalar R. (The curvature scalar only vanishes if the potential has the special form of eq. (3).) For these solutions the overall scale remains undetermined. Only the ratios $\frac{R}{M^2 e^{2S/M}}$, $\frac{<\Phi^\dagger\Phi>}{M^2 e^{2S/M}}$ are fixed and dependent on the parameters of the model.

In presence of gravity it is more convenient to use a version of dilatation symmetry where the metric instead of the coordinates is scaled: $g_{\mu\nu} \to e^{-2\alpha}g_{\mu\nu}$. (Both formulations are related by a general coordinate transformation.) The dilatation current, including the contribution from (5), can be shown to be

$$J^\mu = \sqrt{g}(1 + 12h)Me^{2S/M}\partial^\mu S + \sum_i D_i \chi_i \frac{\partial \mathcal{L}}{\partial(\partial_\mu \chi_i)} \tag{6}$$

In the above χ_i stands for all fields except S and $g_{\mu\nu}$. The conservation of J^μ shows immediately that S is a Goldstone boson [F4], which has only derivative couplings. Indeed we could have made use of rescaled variables $\Phi' = e^{-S/M}\Phi, g'_{\mu\nu} = e^{2S/M}g_{\mu\nu}$ so that only derivatives of S appear in the action.

The above analysis is altered at the quantum level, since dilatation invariance is broken by anomalies [8]. These anomalies just reflect the fact that to renormalize a quantum field theory it is necessary to introduce some scale. The dilatation current (6) is no longer divergenceless so that

$$\partial_\mu J^\mu = \sqrt{g}\Theta^\mu_\mu \tag{7}$$

where Θ^μ_μ is the anomalous trace of the energy momentum tensor. The QCD contribution to this anomalous trace is given by

$$(\Theta^\mu_\mu)_{QCD} = \frac{\beta(g_s)}{2g_s}F_a^{\mu\nu}F^a_{\mu\nu} + m_u\gamma(g_s)\bar{u}u + \cdots \tag{8}$$

whereas the main contribution from weak interactions can be expressed in terms of the weak effective scalar potential $V_{eff}(\Phi, \Phi^\dagger, S)$ [F5]

$$(\Theta^\mu_\mu)_{Weak} = 4V_{eff} - \Phi^\dagger\frac{\partial V_{eff}}{\partial\Phi} - \Phi\frac{\partial V_{eff}}{\partial\Phi^\dagger} - M\frac{\partial V_{eff}}{\partial S} \tag{9}$$

The presence of the additional term in Eq. (7) has two important consequences:

1. The physical laws are no longer invariant under a constant shift in S. In particular, $< \Theta^\mu_\mu >$ will depend on S. As we shall see, possible static solutions correspond to a fixed value $S = S_0$ with $< \Theta^\mu_\mu(S_0) >= 0$.

2. The field S is really only a pseudo Goldstone boson, so that this excitation will acquire some mass .

Before we discuss these points in detail, it is useful to see their parallelism to the invisible axion case. The condition for static S above corresponds to asking that the axion effective potential have an extremum. This extremum condition is what fixes the CP violating parameter $\bar{\theta}$ to zero [4]. The axion, because of the anomaly, also picks up a small mass. Although this is not the way the axion mass is calculated in practice, this mass is related to the second derivative of the axion effective potential at $\bar{\theta} = 0$

The presence of the scale breaking term in Eq. (7) alters the equations of motion for S, adding to it a driving term proportional to Θ^μ_μ. For constant $< \Phi >$ one has from Eq.(6)

$$D^\mu D_\mu e^{2S/M} = \frac{2}{(1+12h)M^2} < \Theta^\mu_\mu(S) > \tag{10}$$

A solution with constant S requires that for some $S = S_0$:

$$< \Theta^\mu_\mu(S_0) >= 0 \tag{11}$$

This solution is stable only if the S field obtains a positive mass term .

$$m_S^2 = -\frac{e^{-2S/M}}{(1+12h)M} < \frac{\partial\Theta^\mu_\mu}{\partial S} >_{S=S_0} \geq 0 \tag{12}$$

We will call the field S a *cosmon* if it fulfils, for some value S_0, the conditions (11) and (12). To have a bearing on the cosmological constant problem, the theory must have another important property, namely that in the vacuum ($S = S_0$) the trace of the energy momentum tensor $< T^\mu_\mu >$ is given by the anomalous trace $< \Theta^\mu_\mu >$. The driving force for the cosmon is then proportional to the cosmological constant . If the field S is somewhere in the vicinity of S_0, it moves due to the nonvanishing driving force. It will perform damped oscillations around S_0 and finally settle at S_0. In this way the cosmological constant is adjusted dynamically to zero [F6]

How does the S dependence of $< \Theta^\mu_\mu >$ arise? If the standard model would be valid up to infinitely short distances, with its dimensionless couplings being defined at some short distance renormalization scale $\bar{\mu}$, the appearance of Λ_{QCD} would be an explicit scale breaking effect. In leading order $\Lambda_{QCD} \simeq \bar{\mu} e^{-b/g_s^2(\bar{\mu})}$. Thus, up to effects arising from the change in the strong β function at fermion mass thresholds, the contribution $< \Theta^\mu_\mu >_{QCD}$ would simply be a (negative) constant $\sim -a(\Lambda_{QCD})^4$. The weak anomaly $< \Theta^\mu_\mu >_{weak}$ is proportional to the fourth power of the expectation value of Φ. If no mass scale $\mu^2 e^{2S/M}$ appears in the Higgs sector and weak symmetry breaking is a radiative effect [16], $< \Phi >$ would again be proportional to $\bar{\mu}$. Therefore also in this case $< \Theta^\mu_\mu >_{weak}$ would be a constant, whose sign depends on the sign of the β-function for the quartic scalar coupling. This in turn depends on the value of the top quark mass. At the other extreme, however, if weak symmetry breaking is essentially due to a (negative) scalar mass term in the Higgs potential, one would have $< \Phi > \sim e^{S/M}$ and therefore $< \Theta^\mu_\mu >_{weak} \sim e^{4S/M}$.

In actual fact we do not believe that the standard model extends to infinitely short distances. Rather we expect that in the neighbourhood of the Planck mass the theory has a larger symmetry: grand unification, higher dimensions and strings may come to play a role. Dilatation symmetry is anomalous if the fundamental dimensionless couplings are running. We may identify the scale of spontaneous breaking of the short distance symmetry with $Me^{S/M}$. Then an additional S dependence of the low energy sector arises through the change of β functions at the scale $Me^{S/M}$. For example, if we define an SU(5) theory by fixing its gauge coupling at a short distance scale $\bar{\mu}$ and if SU(5) is broken to $SU(3) \times SU(2) \times U(1)$ at a scale $\sim Me^{S/M}$, one finds in the one loop approximation, neglecting fermion thresholds,

$$\Lambda_{QCD} \propto M exp[(1 - \frac{\beta_5}{\beta_3})\frac{S}{M}] \tag{13}$$

Here β_5 and β_3 are the β-functions of SU(5) and SU(3) evaluated at the scale $Me^{S/M}$. We conclude that $< \Theta^\mu_\mu >$ has a rather rich and complicated dependence on S, which is sensitive to both long distance and short distance properties of the theory. In view of this, it seems not implausible that the conditions (11) and (12) are fulfilled for some value S_0. The numerical value of S_0 is actually only a matter of convention, since a shift in S can always becompensated by an appropriate multiplicative rescaling of M and all mass parameters in the action. Without loss of generality we choose a convention where $S_0 = 0$, so that μ^2 measures the scalar mass term in the vacuum and M gives the physical scale of high energy spontaneous dilatation symmetry breaking. We have, however, no answer why S settles at a value where $< \Phi >$ is much smaller than M - which is the gauge hierarchy problem.

What about the size of the cosmon mass? Since $< \Theta^\mu_\mu >$ only depends on the dimensionless

combination S/M it follows immediately from Eq.(12)

$$m_S^2 = \frac{m^4}{M^2} \tag{14}$$

where m is a typical scale generated by the dilatation anomalies. There is still much uncertainty about the scale of these anomalies in the weak sector, since it depends on details of the symmetry breaking mechanism. We therefore think it sensible to give only a lower bound, obtained from the presence of the scale Λ_{QCD}

$$m_S^2 \geq \frac{\Lambda_{QCD}^4}{M^2} \tag{15}$$

This gives, for M of order M_P, a limit on the range of the cosmon force [F7]

$$m_S \geq 10^{-11} eV \qquad \lambda \leq 10^4 m. \tag{16}$$

This brings immediately to mind recent speculations about a "fifth" force with similar range [9].

What are the interactions of the cosmon with a macroscopic bulk of matter? For small excitations of S around S_0, one can linearize the field equation (10) for S in the presence of a nucleus N. In contrast to what happens in the vacuum, however, the quantity Θ^μ_μ , when evaluated in a nucleus, is not necessarily the trace of the total energy momentum tensor for the nucleus. The interchange of a cosmon will give rise to a static potential between two nuclei

$$V_{N,N'} = -\frac{G_N}{4\pi}\frac{Q_N Q_{N'}}{r} e^{-\frac{r}{\lambda}} \tag{17}$$

where $G_N = M_P^{-2}$ is Newton's constant. In the above Q_N is the cosmon *charge* of a nucleus, defined by [F8]

$$Q_N = f < N|\Theta^\mu_\mu|N > \tag{18}$$

and where, taking into account the correct normalization of the cosmon kinetic term,

$$f = \sqrt{\frac{16\pi h}{1 + 12h}} \tag{19}$$

Note that apart from the matrix element of Θ^μ_μ between nuclear states - something which we shall be able to estimate reasonably well - Q_N depends on the parameter f entering in Eq.(19). This parameter will serve to characterize the strength of the interactions.

If Θ^μ_μ were the whole trace of the energy momentum tensor, then Q_N would just be proportional to the mass of the nucleus. In this case (18) would just represent an attractive, medium range, modification to the gravitational force. However, Θ^μ_μ is only the anomalous part of T^μ_μ and so one can expect that Q_N will have additional, composition dependent pieces. To get an idea how the composition dependence of the cosmon charge Q_N comes about, we shall make the simplifying assumption that the mass of the nucleus is entirely given by the strongly interacting part of the low energy theory. In practice, this means that both in T^μ_μ and Θ^μ_μ we only retain the QCD pieces. Furthermore, we shall also neglect all contributions from strange and other heavier quarks. With these assumptions one has:

$$Q_N = f < N|\frac{\beta(g_s)}{2g_s} F_a^{\mu\nu} F_{\mu\nu}^a + \gamma(g_s)(m_u \bar{u}u + m_d \bar{d}d)|N > \equiv f < \hat{F}^2 >_N + f_q < \hat{m}_q \bar{q}q >_N \tag{20}$$

while the mass of the nucleus is

$$M_N = < N|\frac{\beta(g_s)}{2g_s}F_a^{\mu\nu}F_{\mu\nu}^a + (1+\gamma(g_s))(m_u\bar{u}u + m_d\bar{d}d)|N > \equiv < \hat{F}^2 >_N + < \hat{m}_q\bar{q}q >_N \quad (21)$$

In the above $\gamma(g_s)$ is the anomalous dimension for the quark operators and

$$f_q = \frac{f\gamma}{1+\gamma} \quad (22)$$

Clearly the difference between (20) and (21) is that the various operator matrix elements enter with different strengths. Indeed, since γ is small, Q_N measures essentially the matrix element of the isosinglet gluon operator in nuclei. Thus, apart from its dependence on M_N, Q_N should depend mostly only on $Z + N = B$. In what follows we shall drop all γ terms, but they can be easily restored if desired.

We may use Eq. (21) to eliminate the gluon operator matrix element in favor of that of the quark operator. This latter matrix element can be estimated from the value of $m_q\bar{q}q$ in nucleons, but it also requires some assumption on what is the contribution of the quark mass terms to the nuclear binding energy. From the measured value of the pion nucleon σ-term, one knows that [10]

$$\sigma = \frac{1}{2}(m_u + m_d) < \bar{u}u + \bar{d}d >_p \simeq 40 - 60MeV \quad (23)$$

while the proton-neutron mass difference, with the electromagnetic contribution extracted, gives [10]

$$\delta = (m_d - m_u) < \bar{u}u - \bar{d}d >_p \simeq 2MeV \quad (24)$$

The corresponding neutron matrix elements follow by an isospin rotation. The nuclear matrix element is obtained by adding the contributions from protons, neutrons and the binding energy:

$$< m_q\bar{q}q >_N = Z < m_q\bar{q}q >_p + N < m_q\bar{q}q >_n - F(\epsilon_B) = (N+Z)\sigma + \frac{1}{2}(N-Z)\delta - x\frac{\sigma}{m_N}\epsilon_B \quad (25)$$

In the above we have introduced a phenomenological parameter, x, which characterizes the possible contribution $F(\epsilon_B)$ of the $m_q\bar{q}q$ operator to the binding energy. If $x = 0$ the binding energy is purely an effect of glue, while if $x = 1$ (neglecting the small δ effect) the $m_q\bar{q}q$ matrix element is proportional to $M_N = m_N(Z+N) - \epsilon_B$.

It is now straightforward to compute Q_N and we find

$$Q_N = f[(1 - x\frac{\sigma}{m_N})M_N - (1-x)\sigma B + \frac{\delta}{2}(Z - N)] \quad (26)$$

If x is small, as we suspect, since nuclear binding should not crucially depend on whether quarks have a mass or not, then Eq. (26) shows an interesting hierarchy of strengths between the mass dependent, baryon number dependent and Z dependent terms [F9]

$$1 : \frac{\sigma}{m_N} : \frac{\delta}{m_N} \simeq 1 : \frac{1}{20} : \frac{1}{500} \quad (27)$$

Although the dominant mass dependent part of Q_N will give rise to an attractive force, the baryon number dependent part will cause the appearance of a composition dependent

repulsive force, as a result of the negative sign in Eq. (26). This sign is just a reflection of the fact that the anomalous part of $< (\Theta^\mu_\mu)_{QCD} >_N$ is dominated by the gluon contribution.

It is useful to compare our expectations with both current experimental knowledge and other theoretical speculations on forces weaker than gravity. For bounds on this new force, where material composition is not important, we need only consider the dominant mass dependent part of the force. Since, to a good approximation, $Bm_N \simeq M_N$, one has

$$Q_N \simeq f(1 - \frac{\sigma}{m_N})M_N \quad (28)$$

Writing the effective potential characterizing possible deviations from Newton's gravity as

$$V(r) = -\frac{G_N M_1 M_2}{r}[1 + \alpha e^{-\frac{r}{\lambda}}] \quad (29)$$

identifies

$$\alpha \simeq \frac{f^2(1 - \frac{\sigma}{m_N})^2}{4\pi} < \frac{1}{3} \quad (30)$$

Obviously, for the force due to cosmons α is positive .In the range below $\lambda \sim 10^3$m ,systematic deviations have been observed between the determination of G_N in mines and on the earth's surface [12]. These discrepancies can be interpreted as evidence for a force weaker than gravity, but which is repulsive with an $\alpha \sim -10^{-2}$. Clearly the cosmon force is in conflict with these observations. However, as de Rújula has pointed out [12], these observations have large systematic uncertainties and are only significant at the level of two times the estimated maximum systematic error. Thus, conservatively, we interpret the mine experiments to place upper bounds $|\alpha| \leq 10^{-2}$ for the range between a few meters and 1 km. For λ larger than 10^3 m bounds on α come from satellite tracking and typically [13] give $|\alpha| \leq 10^{-2} - 10^{-3}$, for λ between 10^3 and 10^4 m.

In experiments sensitive to the composition of the materials studied, then the B and $(N - Z)$ pieces of Q_N will be important. For instance, the difference in force experienced by two test bodies of equal mass M, relative to a third mass $M_\oplus$, is given by the effective potential

$$V_{B,Z}(r) = \frac{G_N M M_\oplus}{r} e^{-\frac{r}{\lambda}}[\alpha_B \Delta(\frac{B}{\mu}) + \alpha_Z \Delta(\frac{Z}{\mu})] \quad (31)$$

Here $\Delta(\frac{B}{\mu})$, $\Delta(\frac{Z}{\mu})$ are the difference in baryon number to atomic mass (in amu) and the difference in proton number to atomic mass (in amu) of the samples, respectively. The strength parameters α_B and α_Z are related to α but reduced (cf. Eq.27)

$$\alpha_B = \alpha\frac{\sigma(1-x)}{m_N} \simeq \frac{1}{20}\alpha \qquad \alpha_Z = -\alpha\frac{\delta}{m_N} \simeq -\frac{1}{500}\alpha \quad (32)$$

One should contrast the result (31) with the suggestion by Fischbach et al.[9] that the anomalies they discovered in their reanalysis of the Eötvös experiment arose from the exchange of a vector excitation which coupled to baryon number [F10]. The potential from such a baryon number dependend force

$$V_{1,2} = \alpha_F \frac{B_1 B_2}{r} e^{-\frac{r}{\lambda}} = \tilde{\alpha}_F \frac{G_N M_1 M_2}{r} (\frac{B}{\mu})_1 (\frac{B}{\mu})_2 e^{-\frac{r}{\lambda}} \quad (33)$$

would also imply the effective potential (31) with $\alpha_Z = 0$ and $\alpha_B = \tilde{\alpha}_F$. However, Eq. (33), implies quite a different material independent residual force. Since $\frac{B}{\mu}$ is very near unity for most materials, Eq.(33) gives an effective potential like that of (29) but with $\alpha = -\tilde{\alpha}_F$. Thus the characteristics of the cosmon force and the, so called, "fifth" force are quite different.

Very recently a group at the University of Washington [14] has carried out a new Eötvös-like experiment and obtained quite strong bounds on α_B, in the range near 10^3 m:

$$|\alpha_B| \leq 2 \times 10^{-4} \qquad 250m \leq \lambda \leq 1400m \tag{34}$$

If we assume the cosmon is in this range, in view of Eq.(32), this implies a bound on the composition independent strength

$$|\alpha| \leq 4 \times 10^{-3} \tag{35}$$

This is a somewhat stronger bound than that available from satellites and mines [13] . It should be pointed out that the Seattle experiment, if a fifth force existed, would place very strong bounds on $\tilde{\alpha}_F$, thereby contradicting the notion that the geological data could be explained by a fifth force.

There is, however, also positive new evidence for a medium range composition dependent force, reported by Thieberger [15]. If his observation is correct, it would imply that at, $\lambda \simeq 10^3$m, $\alpha_B \simeq 3 \times 10^{-3}$, in direct conflict with the Univ. of Washington experiment! This number is also not very comfortable for cosmons, since then α would be quite large ($\alpha \sim 6 \times 10^{-2}$) and be in conflict with the satellite bounds! Of course the cosmon range could be different and then the conflict may ease. It is clear that the present situation is in a very fluid state. Obviously, the correlation between residual torque and baryon number in the Eötvös experiment, noted by Fischbach et al.[9], needs further experimental clarification. We note here only that, if this obtains from a cosmon force, then Eq. (27) explains nicely why the composition dependent part essentially involves only $\Delta(\frac{B}{\mu})$ and not $\Delta(\frac{Z}{\mu})$. Indeed the numbers in Eq. (27) are in agreement with the fit of de Rújula [13] to the Eötvös anomaly in terms of $\Delta(\frac{B}{\mu})$ and $\Delta(\frac{Z}{\mu})$, in which this latter contribution is essentially negligible.

In conclusion, we have presented a model for a dynamical adjustment of the cosmological constant to zero. The main ingredient is dilatation symmetry realized in a Nambu-Goldstone mode and broken only by anomalies. The usual fine tuning of the cosmological constant is replaced, in our approach, by the condition that the trace of the energy momentum tensor in the vacuum is given by its anomalous part $< \Theta^\mu_\mu >$. This condition is not yet well understood. If it holds, however, there is no need any more to carefully cancel the various contributions to the cosmological constant arising from different sectors of the theory. The dynamics selects a value for the cosmon field where this automatically happens. Our model, furthermore,predicts a new force, with a typical range which should be shorter than 10 km, which is mediated by the pseudo Goldstone boson of dilatation symmetry - the cosmon. The couplings of this scalar singlet to ordinary matter are very weak and of strength comparable to gravity. The resulting force for a nucleus has a well understood composition dependence, typified by the parameters which give chiral and isospin breaking in nucleons. The dominant part of the cosmon force is mass dependent and attractive and has to have a strength $\alpha \leq 10^{-2} - 10^{-3}$, so as not to run into conflict with present bounds. It also has a baryon number dependent repulsive component with strength $\alpha_B \sim \frac{1}{20}\alpha$ and a Z dependent component with an even weaker strength. We await eagerly new results searching for weaker forces than gravity.

Acknowledgements

One of us, RDP, acknowledges with pleasure the stimulation of the 1987 Rencontre de Moriond on Exotic Physics. He is also grateful to A. de Rújula for some very useful conversations on the status of fifth forces. J.S. would like to thank F. del Aguila for his interest and for helpful discussions. C.W. would like to thank M. Reuter for useful discussions at an early stage of development of his ideas on this subject. We are grateful to L. Abbott, W. Buchmüller, N. Dragon, J. Ellis, H.P. Nilles, S. Weinberg and A.Zee for useful criticisms of an earlier version of our paper.

Footnotes

[F1] The cosmological constant is known to be really tiny [2]: $< T^\mu_\mu > \leq 10^{-46} GeV^4$

[F2] We speak here of fields very generically, since they need not be the fundamental entities in the theory. Indeed our starting point is much more appropriate for theories without intrinsic scales, like superstrings [6]

[F3] It is well known [7], that in a spontaneously broken theory with dilatation invariance, it is necessary for the scalar potential to develop flat directions in the vacuum, as Eq.(3) does.

[F4] The true Goldstone boson has a small admixture of Φ, which is of order of $< \Phi > /M$.

[F5] This formula follows essentially from the definition of V_{eff} and the role of Θ^μ_μ as a scale breaking effect. If V_{eff} is scale invariant, it has the form $V_{eff} = (\Phi^\dagger\Phi)^4 \mathcal{F}(\frac{\Phi^\dagger\Phi}{M^2 e^{2S/M}})$ and Θ^μ_μ vanishes.

[F6] There have been many other attempts to adjust the cosmological constant dynamically. See, for example, [19].

[F7] We are, of course, aware that for particles with gravitational coupling strength and small mass there can in principle be a cosmological problem, since the energy stored in coherent motion of this excitation is dissipated very slowly. This is similar to the case of the invisible axion [18]. However, this issue depends critically on the initial conditions, which are not obvious in our case. We expect that the expectation value and the effective mass term of the cosmon undergoes important temperature dependent changes (especially at phase transitions where condensates form). Any linear approximation valid for late cosmology becomes invalid for early cosmology (where the initial condition for the late evolution is prepared). Clearly this point deserves further investigation.

[F8] Strictly speaking, the expectation value of Θ^μ_μ in the state stands for the difference between its value in a static nucleus and a piece of vacuum of corresponding volume

[F9] Although, in our approximation, $Z - N$ dependent forces can appear only through the small isospin breaking term δ , in a more complete treatment there will be other isospin breaking contributions in nuclei. They will change the quantitative value of our estimate, but not its magnitude. We note that the dominant part of the cosmon force, not proportional to mass, has a strength proportional to the amount of chiral breaking in nucleons σ/m_N. It is the fact that this breaking is much larger than the effective isospin breaking which gives the baryon number dominance of the residual force!

[F10] Because the strength of the effect in the Eötvös experiment is mostly a function of the nearby mass inhomogeneities, it is not possible to extract a reliable value of $\tilde{\alpha}_F$. A valiant attempt by Talmadge et al.[17] gave $\tilde{\alpha}_F \simeq 10^{-3}$, but this number has an enormous uncertainty.

References

[1] B. Guberina et al., Nucl. Phys. B184 (1981) 476

[2] See for example, S. Bludman and M. Ruderman, Phys. Rev.Lett. 38 (1977)

[3] J. Kim, Phys. Rev. Lett. 43 (1979) 103; M. Shifman, A. Vainshtein and V. Zakharov, Nucl. Phys. B166 (1980) 493; M. Dine, W. Fischler and M. Srednicki, Phys. Lett. 104B (1981) 199

[4] R.D. Peccei and H. Quinn, Phys. Rev. Lett. 38 (1977) 1440, Phys. Rev. D16 (1977) 1791

[5] C.G. Callan, S. Coleman and R. Jackiw, Ann. Phys. (NY) 59 (1970) 42

[6] See for example J. Schwarz in Proc. of the XXIII International Conference in High Energy Physics, Berkeley, July 1986

[7] F. del Aguila and C.D. Coughlan, Phys. Lett. 180B (1986) 25

[8] R. Crewther, Phys. Rev. Lett. 28 (1972) 1421; M.S. Chanowitz and J. Ellis, Phys. Lett. 40B (1972) 397; J.C. Collins, Phys. Rev. D14 (1976) 1965; S.L. Adler, J.C. Collins and A. Duncan, Phys. Rev. D15 (1977) 1712; J.C. Collins, A. Duncan and S.D. Joglekar, Phys. Rev. D16 (1977) 438

[9] A. Dolgov, in The Very Early Universe , eds. G. Gibbons, S.W. Hawking, and S.T. Siklos (Cambridge Univ.Press, Cambridge 1983); F.Wilczek , in Erice Lectures (1983); I. Antoniadis and N. C. Tsamis, Phys. Lett. 144B (1984) 55; L. Abbott , Phys.Lett. 150B(1985) 427; A. Zee , in High Energy Physics, In Honor of P.A.M. Dirac 80th birthday, ed. S.Mintz and A. Perlmutter (Plenum Press); H. P. Nilles, in Proc. of the 9^{th} Johns Hopkins Workshop, Florence 1985, ed. L. Lusanna; S.M. Barr, BNL preprint BNL 38423 (1986).

[10] J. Preskill, M. Wise and F. Wilczek, Phys. Lett. 120B (1983) 127; L. Abbott and P. Sikivie, Phys. Lett. 120B (1983)133; M. Dine and W, Fischler, Phys. Lett. 120B (1983) 137

[11] E. Fischbach et al., Phys. Rev. Lett. 56 (1986) 3

[12] J. Gasser and H. Leutwyler, Phys. Rep. 87C (1982) 77

[13] S.C. Holding et al., Phys. Rev. D33 (1986) 3487

[14] A. de Rújula, CERN preprint CERN TH 4466

[15] A. de Rújula, Nature 333 (1986) 760, Phys. Lett. B180 (1986) 213

[16] C.W. Stubbs et al., Phys. Rev. Lett. 58 (1987) 1070

[17] P. Thieberger, Phys. Rev. Lett. 58 (1987) 1066

[18] S. Coleman and E. Weinberg, Phys. Rev. D7 (1973) 1888

[19] C. Talmadge, S.H. Aronson and E. Fischbach, Proc. of the Rencontre de Moriond 1986