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Pseudo-Goldstone Bosons and New Macroscopic Forces

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Abstract

Pseudoscalar Goldstone bosons may readily be associated with weakly, explicitly broken symmetries giving them mixed CP quantum numbers. In general this leads to *scalar* couplings to nucleons and leptons, which produces coherent long range forces. This can naturally accomodate detectable long range macroscopic forces mediated by bosons completely consistent with conventional cosmological limits, e.g., new interactions within the range of present "5th force" searches which probe a scale of new physics of $f \sim 10^{14}$ GeV.

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While the standard model continues to provide an excellent description of all known phenomena, it leaves so many questions unanswered (e.g. parameters, multiplet structure, Higgs structure, etc.) that many physicists believe there is further new physics beyond the scales currently probed.

A remarkable feature of the standard model is that many of its properties may be derived from the assumption that it is an effective field theory following from an underlying theory relevant at some scale Λ , much larger than the electroweak breaking scale. The states of the standard model are all protected from acquiring the large mass Λ by a symmetry: the vector fields are gauge bosons and the fermions are chiral. The exception is the Higgs scalar, which must either be dynamical, as in technicolor schemes and hence protected by chiral symmetries or must be protected by supersymmetry. The natural question is, if the standard model is indeed an effective theory reflecting underlying physics at a high energy scale, then what are the characteristic signals reflecting this fact? The answer is well known: there should be new terms in the Lagrangian which may give rise to observable phenomena, but, suppressed by inverse powers of the new scale, Λ^{-1} . Hence they give only small effects, difficult to measure unless they give rise to new phenomena beyond the standard model, as is the case for proton decay.

In this letter we wish to discuss another such effect in which a symmetry is realised non-linearly giving rise to a (pseudo) Goldstone mode which plays a role in low energy physics. Its couplings are described by terms in the effective low energy Lagrangian, suppressed by inverse powers of f , the scale of spontaneous breakdown of the symmetry. The phenomenon is well known in the case of the (invisible) axion [1,2] which is employed to suppress the large CP-violating effects of the QCD θ -term. Here an ungauged approximate axial symmetry – the PQ symmetry – (broken in loops only through the $U(1)$ axial-anomaly) is spontaneously broken at a

high scale f_a , thus yielding a pseudoscalar pseudo-Goldstone boson. The coupling to fermions is very feeble, of strength f_a^{-1} , and is a derivative coupling, vanishing as the momentum transfer tends to zero. For large f_a the effects of the axion prove to be very hard to detect experimentally, hence the term “invisible.”

However, it may happen that the non-linearly realised global symmetry is not exact (in the tree approximation) but is an approximate symmetry of the terms of low dimension (≤ 4) in the effective field theory.¹ In this case, the Goldstone mode becomes a pseudo-Goldstone boson [5], does not remain massless (even in the absence of QCD effects) and *can acquire scalar, nonderivative, couplings* in addition to the normal pseudoscalar derivative coupling. This effect is a consequence of the breakdown of CP which occurs when the Goldstone mode acquires a VEV. The particle is spin-zero, but has a mixture of scalar (CP= (+)) and pseudoscalar (CP= (-)) quantum numbers. In view of this schism in the identity of the Goldstone mode, we refer to such a particle as a “schizon.” We emphasize that both the explicit mass of the schizon, as well as its scalar couplings are naturally small since the exact non-linear global symmetry is recovered as they go to zero. The presence of the non-derivative coupling distinguishes the schizon from a conventional pseudo-scalar pseudo-Goldstone boson with pure derivative couplings and justifies the introduction of the new name.

Indeed, one can imagine that the axion itself is a schizon in that the PQ-symmetry is weakly broken at the tree level in addition to its explicit breaking at one loop by the axial anomaly. There will then be a residual θ -angle in QCD, but this will be small, to be consistent with the neutron electric dipole moment, provided that the explicit breaking of PQ-

¹Superstring theories provide an explicit example of this, for after compactification they give rise to new discrete symmetries which may imply new continuous global symmetries for low dimension terms [3]. Indeed, it has been argued that in such theories it is inevitable that the PQ symmetry is explicitly broken [4].

symmetry is small. However, schizons are more general and can exist in addition to the conventional axion. They can thus have masses and decay constants independent of the constraints imposed on those of the axion by its interplay with QCD instantons. The schizon is, however, subject to its own cosmological limits which we consider below.

The scalar coupling component of the schizon is small, suppressed as f^{-1} , but *has no derivatives*, and gives rise to coherent long range forces in matter. There is sufficient freedom in the choice of parameters consistent with cosmological bounds that there may be observable effects, such as reported in recent "fifth force" experiments [6]. These observations may thus signal the presence of structure beyond the standard model characteristic of the new scale f .

We wish first to discuss, in some generality, the phenomenology of schizons. At low energies it is sufficient to consider an effective Lagrangian involving only light fields. For clarity let us first suppress the flavor degrees of freedom, dealing with a single Dirac fermion, ψ . We write the Lagrangian:

$$L_0 = \bar{\psi} i \not{\partial} \psi + (m \bar{\psi}_L \psi_R e^{iq\phi/f} + h.c.) + \frac{1}{2}(\partial\phi)^2 \quad (1)$$

where q is a chiral charge associated with ϕ . L_0 possesses the exact global chiral symmetry,

$$\psi_L \rightarrow e^{iq_L \alpha} \psi_L; \quad \psi_R \rightarrow e^{-iq_R \alpha} \psi_R; \quad q = q_R - q_L; \quad \phi/f \rightarrow \phi/f + \alpha \quad (2)$$

We now allow for explicit breaking of this symmetry by adding terms which break the global symmetry:

$$L = L_0 + (\epsilon \bar{\psi}_L \psi_R e^{iq'\phi/f} + h.c.) + \mu^2 f^2 \cos(\phi/f - \beta) \quad (3)$$

where $q' \neq q$. Note that by shifting ϕ and performing chiral rotations on ψ_L

and ψ_R we can always make m and ϵ real. The parameter β can be viewed as arising from these redefinitions; there is no *a priori* preferred direction for the angle ϕ/f .

Neglecting loops, the effect of the μ^2 breaking term is to force ϕ/f to develop a VEV of $\langle \phi/f \rangle = \beta$, and small fluctuations about this VEV which we call $\hat{\phi}/f = \phi/f - \beta$. The resulting coupling to the field ψ is then given by:

$$(m e^{iq(\hat{\phi}/f + \beta)} + \epsilon e^{iq'(\hat{\phi}/f + \beta)}) \bar{\psi}_L \psi_R + h.c. = \rho(\hat{\phi}) e^{i\sigma(\hat{\phi})} \bar{\psi}_L \psi_R + h.c. \quad (4)$$

where we find trivially:

$$\rho^2 = m^2 + \epsilon^2 + 2m\epsilon \cos[(q - q')(\hat{\phi}/f + \beta)] \quad (5)$$

$$\sin(\sigma) = \rho^{-1} \{ m \sin[q(\hat{\phi}/f + \beta)] + \epsilon \sin[q'(\hat{\phi}/f + \beta)] \} \quad (6)$$

Now, we may perform the local redefinition of the fields, $\psi_L \rightarrow e^{i\sigma/2} \psi_L$; $\psi_R \rightarrow e^{-i\sigma/2} \psi_R$, to obtain:

$$L = \bar{\psi} i \not{\partial} \psi - \frac{1}{2} \partial^\mu (\sigma) \bar{\psi} \gamma^5 \gamma_\mu \psi + \rho \bar{\psi} \psi + \frac{1}{2} (\partial \hat{\phi})^2 + \mu^2 f^2 \cos(\hat{\phi}/f) \quad (7)$$

so, the phase $\sigma(\hat{\phi}) \sim q(\hat{\phi}/f + \beta)$ behaves like the conventional axion, and is derivatively coupled.

However, we see that the theory now has a net scalar-Yukawa coupling of $\hat{\phi}$ to $\bar{\psi}\psi$ contained in ρ . Throughout we shall assume that $\epsilon \ll m$, and the scalar coupling term is of the form:

$$L_Y \approx \left[\frac{(q' - q)m\epsilon \sin[(q - q')\beta]}{(\sqrt{m^2 + \epsilon^2 + 2m\epsilon \cos[(q - q')\beta]})} \right] \frac{\hat{\phi} \bar{\psi} \psi}{f} \quad (8)$$

Notice that there would be no such term in the limit $\epsilon = 0$, even though $\mu^2 \neq 0$; both ingredients are required to produce both the scalar mass and

the scalar-Yukawa coupling. The nontrivial minimum of the potential in eq.(3) forces the ϕ field to develop a VEV, which violates CP. In this simple model, the overall CP-phase, σ , could be rotated away. The manifestation of the CP-violation is the mixture of pseudoscalar and scalar couplings of ϕ .

With an arbitrarily small mass, μ , the exchange of the schizon can give rise to an intermediate range classical force in bulk matter, $\sim e^{-\mu r}/r^2$. The strength of this force between masses M_a and M_b composed of the ψ particles is $\sim G' M_a M_b$ where the constant G' is, relative to Newton's gravitational constant, given by

$$\frac{G'}{G_N} = \frac{\epsilon^2 (q - q')^2 M_{Planck}^2 \sin^2[(q - q')\beta]}{m^2 f^2} \quad (9)$$

This result forms the basis of the present paper.

Of course, we have presently neglected the effects of loops which will produce corrections to the schizon mass term of eq.(3). The leading perturbatively induced schizon mass is contained in the amplitudes of the form, $\langle T\rho(\hat{\phi})\bar{\psi}\psi \rangle$. These are superficially quadratically divergent, which would be a disaster from the point of view of controlling the smallness of the breaking. This problem can be cured by introducing additional fermions and/or symmetries into the coupling of the pseudoscalar in eq.(1), and is expected in the context of realistic theories. We denote by M^2 (M is quite distinct from f and we expect generally that $M \sim M_W$), the characteristic (mass)² scale of these effects. For example, we can appeal to supersymmetry above M_W to cancel the quadratic divergence in the schizon mass, whence we find generically that $M^2 \sim M_{SUSY}^2$. It is also possible to arrange a GIM cancellation of the quadratic divergence amongst the existing generations of quarks and leptons whence $M^2 \sim m_{charm}^2$. In our effective field theory, below the scale of such physics, the radiative effects should be interpreted as real effects which may be used to bound M^2 . The induced schizon mass

term is then $\sim [m\epsilon M^2/8\pi^2] \cos[(q - q')(\hat{\phi}/f + \beta)]$, up to logarithms.

There can also arise an induced mass term from the breaking of chiral symmetry contained in the amplitude $\langle T\rho(\hat{\phi})\bar{\psi}\psi \rangle \sim \rho\Lambda_{QCD}^3$. In this case we induce a mass term for the schizon of $\sim [\epsilon\Lambda_{QCD}^3] \cos[(q - q')(\hat{\phi}/f + \beta)]$. Thus, we see that the general form for the induced schizon mass is given by $[\epsilon K^3] \cos[(q - q')(\hat{\phi}/f + \beta)]$ where K^3 is the larger of $\sim mM^2/8\pi^2$ or Λ_{QCD}^3 . The effect of including these additional terms in the potential for the schizon is an additional shift in the VEV, i.e., $\hat{\phi}$ now acquires a VEV.

For the sake of convenience below, we shall group together these various contributions to the potential of the schizon associated with explicit breaking by defining:

$$P(\phi/f) = -\mu^2 f^2 \cos[\phi/f - \beta] - [\epsilon K^3] \cos[(q - q')(\phi/f)] + \dots \quad (10)$$

and we shall define the minimum of $P(\phi/f)$ to occur for $\phi/f = \delta$. About this minimum we have the small fluctuations,

$$P(\hat{\phi}/f + \delta) \approx \frac{1}{2} m_\phi^2 \hat{\phi}^2. \quad (11)$$

Let us now consider the modifications necessary to include the flavor degrees of freedom of the standard model. For the moment we consider only possible coupling to quarks and we must face the problem of maintaining at least an approximate PQ-symmetry to cancel the QCD θ -term. We consider first the case of an *exact* PQ-symmetry with an axion and an independent schizon field. The generalization to the light quarks of eqs.(1, 3) is a coupling to fermions in the form:

$$\begin{aligned} L_0 = & m_u \bar{u}_L u_R e^{ia/f_a} + m_d \bar{d}_L d_R e^{ia/f_a} + h.c. \\ & + \epsilon_u \bar{u}_L u_R e^{i(q_u \phi/f_\phi + a/f_a)} + \epsilon_d \bar{d}_L d_R e^{i(q_d \phi/f_\phi + a/f_a)} + \theta G_{\mu\nu} \tilde{G}^{\mu\nu} \end{aligned} \quad (12)$$

where we exhibit also the QCD θ -term. Of course, we could include a series of terms such as the ϵ_q , but the above Lagrangian is sufficient for illustration.

We may write for the quark mass terms:

$$m_q e^{ia/f_a} + \epsilon_q e^{i(q_q \phi/f_\phi + a/f_a)} = \rho_q e^{i(\sigma_q + a/f_a)} \quad (13)$$

and perform chiral rotations to bring L_0 into the form:

$$\begin{aligned} L_0 = & \rho_u \bar{u}_L u_R + \rho_d \bar{d}_L d_R + h.c. + \mu^2 f_\phi^2 \cos(\phi/f - \beta) \\ & + [\theta - c(\sigma_u + \sigma_d + 2a/f_a)] G_{\mu\nu} \tilde{G}^{\mu\nu} \end{aligned} \quad (14)$$

where $c = g_{QCD}^2/32\pi^2$.

The effect of instantons is to break the PQ-symmetry by producing an effective potential which is an even function of the coefficient of the $G\tilde{G}$ term in eq.(14). Similarly, we have the effects of quark loops and chiral symmetry breaking which induce mass terms for the schizon, in addition to the explicit terms as discussed above. With $\epsilon_q \gg m_q$, the resulting effective Lagrangian for the schizon and axion takes the form:

$$\begin{aligned} L_{a,\phi} = & \frac{1}{2}(\partial a)^2 + \frac{1}{2}(\partial \phi)^2 - P(\phi/f_\phi) \\ & - \frac{1}{2} m_a^2 f_a^2 (\theta/2c - \sigma_u/2 - \sigma_d/2 - a/f_a)^2 \end{aligned} \quad (15)$$

The schizon may be written in the shifted form $\phi/f_\phi = \hat{\phi}/f_\phi + \delta$ to minimize $P(\phi/f_\phi)$. Then, in our approximation, $\sigma_q \rightarrow (q_q \epsilon_q/m_q)(\hat{\phi}/f_\phi + \delta)$. The axion will then be shifted to cancel $\theta/2c - \sum_q (q_q \epsilon_q/m_q)(\delta/2)$. The result of this is to zero the constant coefficient of $G\tilde{G}$, which thus completely eliminates the strong CP problem, and $L_{a,\phi}$ then takes the form:

$$L_{a,\phi} = \frac{1}{2}(\partial \hat{a})^2 + \frac{1}{2}(\partial \hat{\phi})^2 - \frac{1}{2} m_a^2 f_a^2 [(q_u \epsilon_u/m_u + q_d \epsilon_d/m_d) \hat{\phi}/f_\phi + \hat{a}/f_a]^2$$

$$-\frac{1}{2}m_\phi^2\hat{\phi}^2 \quad (16)$$

Here m_ϕ^2 is the resulting mass of the small fluctuation and is of order $\text{Max}(\mu^2, \sum_q q_q^2 \epsilon_q K_q^3 / f_\phi^2)$

The overall θ -term is cancelled in the usual way, since our model maintains the PQ-symmetry, however there remains mixing between the schizon and the axion. We will be interested in the range of parameters in which $10^9 < f_a < 10^{12} \text{Gev}$, for consistency of the axion with conventional astrophysical limits [7,8], and $f_\phi \gg f_a$, $m_\phi^2 \ll m_a^2$ (and $\epsilon_q/m_q \ll 1$), for an interesting range and strength of the schizon mediated force. In this case, one readily finds that the (mass)² eigenvalues of eq.(16) are $m_a^2(1 + O(q_q^2 f_a^2 \epsilon_q^2 / f_\phi^2 m_q^2))$ and m_ϕ^2 . The axion will develop a negligible scalar coupling in this limit through its mixing with ϕ , of order $O(f_a \epsilon_q / f_\phi m_q)$, and, conversely, the schizon develops a further negligible correction in its coupling to the anomalies.

The resulting scalar Yukawa couplings of the schizon is contained in the terms:

$$\begin{aligned} L_0 &= \rho_u \bar{u}_L u_R + \rho_d \bar{d}_L d_R + h.c. \\ &\rightarrow m_u \bar{u}_L u_R + m_d \bar{d}_L d_R + \epsilon_u \sin(\delta) \frac{q_u \hat{\phi} \bar{u}_L u_R}{f_\phi} \\ &\quad + \epsilon_d \sin(\delta) \frac{q_d \hat{\phi} \bar{d}_L d_R}{f_\phi} + h.c. + \dots \end{aligned} \quad (17)$$

and, again, arises essentially as a consequence of the offset angle, δ , reflecting the presence of CP-violation.

At this point several remarks are in order. It is straightforward to extend eq.(12) to additional generations and it is interesting that, in a two generation case, we may choose a GIM cancellation of the quadratic divergences, as discussed above, by choosing, e.g., $\epsilon_u = \eta m_c$, and $\epsilon_c =$

$-\eta m_u$, etc., where $\eta \ll 1$. Such a model can readily be generated from a realistic Higgs potential, but we can also appeal to SUSY to eliminate the quadratic divergences. Note, furthermore, that by choosing $\epsilon_u = -\epsilon_d$, we eliminate the chiral condensate contribution to the schizon mass and we would have a coupling of the schizon to I_3 . Also, we see that the chiral rotation of eq.(14) will induce a coupling of the schizon, in addition to the axion, to the electromagnetic anomaly, $\frac{\alpha}{4\pi}(\frac{4}{9}\sigma_u + \frac{1}{9}\sigma_d)F\tilde{F}$, hence schizons can in principal be produced in $E \cdot B$ experiments at a suppressed rate, but at considerably longer wavelengths than those usually assumed for the axion for the low mass schizon of interest to us here. Also, we've considered only flavor diagonal couplings; with off-diagonal couplings the possibility of CP-violation in strong classical schizon fields arises. Conceivably the terrestrial schizon field plays a role in $K_L K_S$ CP-violation, and measurements above the surface of the earth might show variations.

Notice that the relative signs and strengths of ϵ_u and ϵ_d are arbitrary. Thus, the schizon can couple to a linear combination of B and I_3 (and the leptonic generalization of this result is an arbitrary linear combination of N_e and N_{ν_i}). This contrasts with the cosmon model which has difficulty producing macroscopic force coupled to I_3 [9]; in principle, with $\epsilon_u = -\epsilon_d$, we can have pure coupling to I_3 .

We turn now to the cosmological limits on the parameters describing the schizon. First, it is readily seen that the conditions necessary to maintain thermal equilibrium of the schizon in the early Universe are generally not met. As seen in eq.(17), the coupling strength of the schizon to massive states is characterized by the coupling strength $\sim \epsilon/f$ where typically $\epsilon \ll m_{Baryon}$ (though heavier quarks may enter). As the temperature of the Universe falls below $\sim f$ the massive states decouple and the schizon interaction becomes weak. In order to maintain thermal equilibrium we require that typical reaction rates exceed the Hubble expansion

rate, $n\sigma \gg H \sim T^2/M_P$ where $\sigma \sim \epsilon^2/f^2 T^2$, and $n \sim T^3$. This only applies provided $T > m_{Baryon}$ and we thus will not remain in thermal equilibrium at any subsequent time provided:

$$m_{Baryon} < \left(\frac{m_{Baryon}^2 f_\phi}{\epsilon^2 M_P} \right) f_\phi \quad (18)$$

Since this condition is generally expected to be satisfied, the schizon evolution will be governed by its equation of motion:

$$\ddot{\phi} + 3H\dot{\phi} + m_\phi^2 \phi = 0 \quad (19)$$

For $H > m_\phi$ the schizon field is critically damped, thus, effectively frozen to a constant value and its potential energy of order $m_\phi^2 \phi^2 \sim m_\phi^2 f^2$ does not change.

The ensuing cosmological behavior depends critically upon the terms giving rise to m_ϕ . The terms arising at the scale f_ϕ and the fermion loop contributions are present at all temperatures and essentially constant. We shall refer to this as the $m_{\phi 0}$ contribution to the ϕ mass. The nonperturbative contribution involves the chiral condensate, $m_{\phi 1}^2 \sim \epsilon \langle \bar{q}q \rangle / f$, and appears only below a temperature $T_1 \sim O(\Lambda_{QCD})$. Once H drops below $m_{\phi 0}$ at a temperature $T_0 \sim \sqrt{m_{\phi 0} M_P}$, then ϕ freely oscillates and redshifts as a matter field with energy $\approx R^{-3}$:

$$\rho_\phi(T_1) = \rho_\phi(T_0) \frac{T_1^3}{T_0^3} = \frac{1}{2} m_{\phi 0}^2 f_\phi^2 \frac{T_1^3}{T_0^3} \quad (20)$$

As the temperature drops below T_1 , $m_{\phi 1}^2$ switches on, shifting the minimum of $P(\phi/f_\phi)$ and releasing further latent heat $\sim \frac{1}{2} m_{\phi 1}^2 f_\phi^2 \alpha$ where: $\alpha \approx \left[\frac{m_{\phi 1}}{m_{\phi 0}} \right]^2 \leq 1$. Adding this to $\rho_\phi(T_1)$ and evolving to lower tempera-

tures gives, for $T < T_D$:

$$\frac{\rho_\phi(T)}{\rho(T)} = \frac{1}{2} \frac{f_\phi^2}{M_P^2} \frac{T_0}{T_D} \left[1 + \alpha^2 \left(\frac{T_0}{\text{Min}(T_1, T_0)} \right)^3 \right] \leq \frac{\Omega_\phi}{\Omega_{\text{visible}}} \quad (21)$$

where $\rho(T)$ is the total visible energy density of the Universe, Ω_ϕ ($\Omega_{\text{visible}} \sim .1$) is the fraction of closure density of the coherent oscillations of the schizon (visible matter) and T_D the temperature at transition to a matter dominated Universe, $T_D \sim 10\text{eV}$ (i.e., T_D is defined by $\rho_{\text{matter}} = \rho_{\text{rad}}$). Saturating eq.(21) with $\Omega_\phi = 1$ corresponds to assuming that the schizon coherent oscillations close the Universe and thus account for the dark matter.

A second astrophysical limit comes from the cooling rates of Red Giant stars [8]. In this case schizon cooling proceeds via the electromagnetic anomaly only, and for the interesting range $\epsilon/f_\phi \ll m_q/f_a$, cooling due to schizon radiation is negligible compared to that of the axion, and plays no significant role in the further discussion.

With the above cosmological limits, we may now discuss the generic long-range forces that may arise in bulk matter due to schizon exchange. The coherent long range force resulting from the schizon scalar interaction has range m_ϕ^{-1} and strength given by:

$$\frac{G'}{G_N} \approx \left(\sum \frac{\epsilon_q q_q M_P \sin \delta}{m_N f_\phi} \right)^2 \leq \left(\sum \frac{m_\phi^2 f_\phi M_P \tan \delta}{K_q^3 q_q^2 m_N} \right)^2 \quad (22)$$

where m_N is a nucleon mass and we have used the bound on radiative or non-perturbative masses discussed above (for consistency we must check $\epsilon_q < m_q$). Combining this result with the cosmological bound of eq.(21) gives:

$$\frac{G'}{G_N} \leq \left(\frac{m_\phi \cdot (10^{17})}{\text{GeV}} \right)^{(7/2)} K_q^{-6} \left[1 + \alpha^2 \left(\frac{T_0}{\text{Min}(T_1, T_0)} \right)^3 \right]^{-1} \quad (23)$$

where we have taken $\tan \delta \approx q_q \approx 1$ as reasonable values. This equation shows that macroscopic forces of range and strength measureable in recent experiments [6] may result from schizon exchange, consistent with the cosmological bounds. For example, $m_\phi \approx 10^{-17} \text{ GeV}$ gives a range of 500 meters, $T_0 \sim 10 \text{ GeV}$, and using eq.(21) $f_\phi \leq 10^{14} \text{ GeV}$. Using a large cutoff, $M^2 = O(100 \text{ GeV})^2$, as in a SUSY theory, we obtain $K \sim 1$; $\alpha \sim O(10^{-3})$, and thus from eq.(23) $G'/G_N \leq 1$. Using a small cutoff, $M^2 = O(1 \text{ GeV})^2$ as in GIM type cancellation, gives $K \sim 10^{-1}$; $\alpha \sim 1$, and, again $G'/G_N \leq 1$. In both cases $\epsilon_q \ll m_q$, so the bound of eq.(23) applies. The upper limit on the strength of the new force should be compared with the current experimental results, interpreted as bounds, $G'/G_N \leq 10^{-2}$ [10], showing that measureable macroscopic forces may easily arise via exchange of a schizon which has parameters fully consistent with standard cosmology.

We note that another suggestion for a long range force mediated by a "cosmon" [9], the pseudo-Goldstone boson associated with the dilatation symmetry, could not satisfy the constraints of cosmology if its properties are such as to give a force in this range [11].

There are many other possible couplings of the schizon. It may couple to some combination of lepton number alone, in which case there will be no mixing with the axion. In this case there will be no contribution to the schizon mass from a fermionic chiral condensate, but there will remain radiative mass generation. As a result the bounds on the strength relative to range remain as given in eqs.(23, 24) without the α^2 term. In this case there is also a constraint from Red Giant cooling as the schizon now couples to the electron plasma directly and $\epsilon_e q_e \cos(\delta)/f_\phi \leq 1.4 \times 10^{-13}$.

At the other extreme, we may identify the schizon with the axion itself, allowing for explicit breaking of PQ-symmetry. In this case we identify a with ϕ in eq.(12). This leads to eq.(15) with $a = \phi$, but in this case the term $P(a/f_a)$ will force the axion to develop a VEV which, now, will not exactly

cancel the constant coefficient of $G\tilde{G}$, since the PQ-symmetry is explicitly broken by the terms giving rise to $P(a/f_a)$. To maintain the cancellation of θ to the required precision, of order $\leq 10^{-9}$ [12], the term $P(a/f_a) \approx \frac{1}{2}\mu^2(a-\delta)^2$ must be much smaller than in the preceding case, $\mu^2 \leq 10^{-9}m_a^2$. The coherent long range force resulting from schizo-axion exchange has range m_a^{-1} and strength given by the analogue of eq.(22) modified to take account of the stronger mass bound which is given by eq.(22) with $m_\phi^2 \rightarrow 10^{-9}m_\phi^2$ on the right hand side. This time, however, the cosmological bounds are $10^{10} < f_a < 10^{11}$ [7,8] and $m_a \approx \Lambda_{QCD}^2/f_a$ is essentially known. Thus, with reasonable estimates for the other parameters we have:

$$\frac{G'_a}{G_N} \approx 10^{-4} \text{ to } 10^{-10} \quad (24)$$

with range $\approx 10^{-3}$ to 10^{-4} meters.

In conclusion, we have shown that the standard model, with reasonable and minimal extension, may readily accomodate a wide range of previously uninvestigated phenomena which are observable and which would foreshadow the existence of physics at a new scale, $f \gg M_W$. The essential result is that there may exist new macroscopic forces with range and strength in the region probed by recent "5th force" experiments. These forces are mediated by pseudo-Goldstone bosons with scalar couplings induced by explicit breaking of the associated symmetry. The properties of these bosons, which we have dubbed "schizons," are consistent with standard Big Bang cosmology provided the strength is bounded by $\approx [(10 \text{ to } 100 \text{ meters})/\lambda]^{7/2} \times (\text{gravitational})$, for a range of the force $\lambda = m_\phi^{-1}$. Coherent relic schizon fluctuations may even account for the dark matter in the Universe.

Schizons may occur in any model with spontaneously broken discrete symmetries which give rise to global symmetries when the effective Lagrangian is restricted to terms of low dimension. The coupling to matter is

not determined by the low energy theory – thus coupling to conventional matter may involve an arbitrary linear combination of Baryon number, B , strong isotopic spin, I_3 , or lepton numbers, L_i and L_{ν_i} . This can give rise to composition dependent long range forces of the type probed by Eötvös-type experiments. Being a scalar mediated force, however, it will not produce a repulsive force between like sign charges and hence is unlikely to explain the reported 2σ effect seen in geophysical measurements. It should be noted that most of the higher precision experiments do not address the repulsive/attractive nature of the force [6]. Since schizons violate CP, they may lead to other macroscopic effects as well [13].

Schizons provide one of the few windows on the high energy world beyond the standard model, one we believe to be of equal importance with other windows such as proton decay. Current experiments sensitive to schizon-mediated forces are capable of probing new scales of fundamental physics, e.g., for a range of 500 meters we see from eq.(21) that the scale probed is up to $f \sim 10^{14}\text{GeV}$! Our analysis has shown there is a wide range of parameters for such forces consistent with the standard model and cosmology. The phenomenology and more explicit models of schizons will be discussed in greater detail elsewhere [14].

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