Turning on gravity with the Higgs mechanism

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25 February 2016

Abstract. We investigate how a Higgs mechanism could be responsible for the emergence of gravity in extensions of Einstein theory, with a suitable low energy limit. In this scenario, at high energies, symmetry restoration could "turn off" gravity, with dramatic implications for cosmology and quantum gravity. The sense in which gravity is muted depends on the details of the implementation. In the most extreme case gravity's dynamical degrees of freedom would only be unleashed after the Higgs field acquires a non-trivial vacuum expectation value, with gravity reduced to a topological field theory in the symmetric phase. We might also identify the Higgs and the Brans-Dicke fields in such a way that in the unbroken phase Newton's constant vanishes, decoupling matter and gravity. We discuss the broad implications of these scenarios.

Keywords Cosmology, quantum gravity, scalar-tensor theories, Higgs mechanism

1. Introduction

The Higgs mechanism is a central aspect of modern particle physics [1, 2, 3]. At its core lies the idea that the vacuum and its perturbative excitations do not partake of the full set of symmetries of the underlying theory. This happens because the Higgs field acquires a non-trivial vacuum expectation value (VEV), and fluctuations around it are not as symmetrical as the Lagrangian of the theory. The symmetries thus spontaneously broken are expected to be explicitly restored at high energies, when the potential loses its 'Mexican hat' shape, allowing the field to rest at the origin. The Higgs mechanism is responsible for giving masses to gauge bosons (which must be fundamentally massless, due to the gauge symmetry), as well as to fermions.

The possibility that the Higgs mechanism might have gravitational content is intriguing, and has been considered before both in a theoretical framework, with very different slants, all orthogonal to the approach in our Letter (e.g. [4, 5, 6]), and from the observational point of view, e.g. regarding spectral lines [7]. In this Letter we consider a radical prospect in this regard: that gravity only properly "switches on" once the Higgs field acquires a non-trivial VEV. This could occur in relation to the electroweak Higgs doublet or with any other Higgs-like field, associated with any group or pattern of symmetry breaking. In any such scenario, as we probe higher and higher energies, the symmetries spontaneously broken by the Higgs field are eventually restored, at which point gravity "turns off", dropping out of the dynamical picture to some degree.

The precise sense in which this turn-off happens depends on the technicalities of the implementation of this broad idea. For example, we may adopt the view that Einstein gravity results from applying a constraint (the so-called simplicity constraint) to a topological field theory (the "BF" theory introduced or reviewed in [8, 9, 15], and defined in the next Section). The constraint enters the action via a Lagrange multiplier, but suppose that the modulus of the Higgs field multiplies this term. Then, in the symmetric phase (with $\phi_0 = 0$), the constraint would be turned off, rendering gravity a topological field theory, with the dynamical degrees of freedom unleashed only in the broken phase (with $\phi_0 \neq 0$, turning on the constraint). In this realization, at high energies there are no gravitons. In addition the vacuum solutions are all of a topological nature. For another discussion of the possible gravitational relevance of the Higgs mechanism see ref [16], and for a different approach see ref [17].

Another possibility follows from identifying the Higgs field (or any Higgs-like field) with the Brans-Dicke field [22] or any variation thereof in a more general scalar-tensor gravity theory [23], so that the Higgs VEV determines the strength of the gravitational interaction, as fixed by the effective Newton's "constant" G_{eff}. By choosing a suitable coupling function between the field and the Ricci scalar, we can ensure that $G_{eff} = 0$ in the unbroken phase, but retain $G_{eff} \neq 0$ in the broken phase. This effectively decouples matter and gravity in the unbroken phase, so that there are still free gravitons at high energies, but there is no way to produce them other than by vacuum quantum fluctuations. The self-interaction of the graviton is also turned off. In this scenarios there are still non-trivial vacuum solutions, just as with the usual scalar-tensor theories; however singularities induced by matter are no longer present.

In our model we make a minimum of theoretical assumptions. In particular, we will only assume the standard model and general relativity, or similar gravity theory. These provide gauge and general covariances. By combining general covariance with the SU(2) gauge invariance, the Higgs field is allowed to couple with general relativity with an SU(2) invariant (singlet representation) of a functional of the Higgs field, $f(\phi)$ ‡. So, all gauge invariant functions are allowed including Higgs multiplets in our model, as in any gauge invariant extension to the standard model. The only difference is that the gauge invariant functional couples to the Ricci scalar. Even with a Higgs multiplet, it is important that the invariant functional, $f(\phi)$, vanishes in the symmetric phase. So, in general, if one component of the multiplet is vanishing, it will not be sufficient just to satisfy $f(\phi) = 0$ to switch gravity off. In the Higgs mechanism, the mass terms arise from the Higgs VEV, and this is where gravity turns off. Our model assumes the minimal Higgs doublet, so there is only a VEV for the Higgs. The various values of the Yukawa couplings, which control the masses of the matter fields, do not play a role in switching gravity on or off—only the Higgs VEV controls that.

2. The topological model

Let us consider in more detail the first of these possibilities. Consider the modified BF action [8, 9, 15]:

$$S_{BF} = \int B^{IJ} \wedge F_{IJ} - \frac{1}{2}|\phi|\Phi_{IJKL}B^{IJ} \wedge B^{KL}, \tag{1}$$

where $F^{IJ}(\omega)$ is the curvature of the connection ω^{IJ} taking values in an algebra (assumed here to be $SO(3,1)$ or $SO(4)$, but it could also be $SU(2)$); B^{IJ} is a 2-form taking values in the same algebra, and Φ_{IJKL} is a Lagrange multiplier enforcing the simplicity constraint. Here, the indices run as $I = \{0, i\} = 0, 1, 2, 3$, $\mu = 0, 1, 2, 3$. If $|\phi|$ is a non-vanishing constant it is known that this is a constrained version of BF theory such that the action is equivalent to the Cartan formulation of Einstein theory (i.e. a first-order formulation which reduces to the usual second order in the absence of spinors; see [10, 11] for reviews). However, if $|\phi| = 0$ (i.e. if the simplicity constraint is not enforced) the theory is just plain BF theory, i.e. it is a topological field theory. In both cases a densitized metric can be constructed according to:

$$\sqrt{-g}g_{\mu\nu} = \frac{1}{12}\epsilon_{ijk}\,\epsilon^{\alpha\beta\gamma\delta}B^i_{\mu\alpha}B^j_{\beta\gamma}B^k_{\delta\nu}. \tag{2}$$

and let us assume a convention so that a Lorentzian metric has signature $- + ++$. The field ϕ can be a real scalar, or a complex doublet, depending on how close to the

‡ The function $f(\phi)$ can in principle arise from string theory due to the compactification (dimensional reduction) of the full 10 dimensional theory to our 4 dimensional world. String theory will yield a volume factor containing information of the extra six dimensions in the form of a function that could in principle have the required features of our function. This calculation is beyond the scope of this letter and we shall pursue it in an upcoming work

standard model we want the model to be. All we require is that ϕ be Higgs-like, so that its action is of the form:

$$S_{Higgs} = \int d^4x \sqrt{-g} \left(-\frac{1}{2}(\partial_\mu \phi)(\partial^\mu \phi) - V(\phi) \right) \tag{3}$$

where $V(\phi)$ is a 'Mexican hat' potential, which at high temperatures is flattened into a simple bowl.

This theory has two types of solutions corresponding to two phases, connected by a phase transition. At low temperatures, $V(\phi)$ has a 'Mexican hat' shape, so that the vacuum solutions lie on its rim, with the non-trivial VEV $|\phi| = \phi_0 \neq 0$. Even though there is a topological sector [20, 21] in this phase, there is also a dynamical sector, following from the solution:

$$B^{IJ} = \frac{1}{2}\epsilon^{IJ}{}_{LM} e^L \wedge e^M \tag{4}$$

written in terms of e^I (to be interpreted as the tetrad). For non-degenerate tetrads, this can be inserted back into action (1), yielding the Einstein-Cartan action. The torsion-free solutions then reduce this action to the Einstein-Hilbert action. At high temperatures, however, $V(\phi)$ becomes a bowl potential, with a vacuum solution at $\phi = 0$, so that the simplicity constraint cannot be imposed. Then, as is well known, only the topological sector exists.

The implications for cosmology and quantum gravity are far-reaching. It would appear that the Big Bang model at high temperatures gives way to a simple topological solution [12]. For isotropic cosmologies where the spatial curvature is zero the cosmological singularity is therefore removed§. This is indeed the case in this model, where at high energies the Higgs field switches off the graviton degrees of freedom. The quantization of the ensuing topological field theory is standard [27].

A further layer can be uncovered by allowing the Φ_{IJKL} to have a trace (a standard way to introduce a cosmological constant). Symmetry dictates that we can have a renormalizable coupling of a Higgs-dependent cosmological constant to gravity, $\Lambda_0 \to \Lambda(\phi)$. It turns out that this leads to even more interesting implications. The action then becomes:

$$S_{BF\Lambda} = S_{BF} + \Lambda(\phi) B^{IJ} \wedge B_{IJ}. \tag{5}$$

Here we identify the cosmological constant with the Higgs potential $\Lambda(\phi) = V_H(\phi) = \lambda(\phi^2 - \eta^2)^2$. Upon variation with respect to the Higgs we get the following condition

$$D_\mu D^\mu \phi = \Phi_{IJKL} B^{IJ} \wedge B^{KL} + \Lambda(\phi)' B^{IJ} \wedge B_{IJ} \tag{6}$$

and the Einstein equations,

$$F^{IJ} = B^{IJ} + \Lambda(\phi)' B^{IJ} + L_\phi B^{IJ}. \tag{7}$$

Here L_ϕ is the Higgs field Lagrangian. In the symmetric phase, $\phi = 0$, and $\Lambda = \Lambda_0$, so that $F^{IJ} = B^{IJ}$, which yields the equations of topological gravity with a non-vanishing

§ An anisotropic solution could still have a singularity, for example becoming Kasnerlike as in the non-topological phase. We intend to pursue this issue in a forthcoming work cosmological constant. These theories have been dealt with in Lorentzian spin-foam quantization and also in relation with the Kodama state [13, 14, 18, 19].

Here, we see that the missing link regarding the issues of non-normalizability in such theories may be the Higgs dependence of the cosmological constant. The non-normalizability of topological quantum gravity with a cosmological constant may be signifying an instability to decay to a state of vanishing cosmological constant. For example, semi-classically in the symmetry-breaking phase, $\phi = \phi_{min}$, the cosmological constant vanishes and we get dynamical gravity. There is an intermediate state where the tachyonic instability sets in and the Higgs is rolling down to its minimum. Here we see that the cosmological constant is dynamical and we have a situation where gravity is turned on and the cosmological constant is varying with the evolution with the Higgs field. The Higgs can tunnel from the false vacuum to the true vacuum and this may serve as a potential way to explain why the true cosmological constant might be zero, or at least much smaller than $\mathcal{O}(10^{-122})$, today. This would require the dark energy responsible for the acceleration of the universe to be supplied by a slowly varying energy field that has come to dominate the universe at late times.

3. The scalar-tensor varying-G model

As a second, less radical realization for the proposal in this Letter, we could seek to identify the Brans-Dicke field, or similar, and a Higgs-like scalar field undergoing symmetry restoration at a phase transition [4]. However, such a generic model must be fundamentally modified before it is put to service regarding the main idea in this paper: "switching off" gravity in the unbroken phase. We seek a model in which $G_{eff} = 0$ in the unbroken phase, when $\phi = 0$. But in the simplest such model one has $G_{eff} \sim 1/\phi$, so that $G_{eff} \to \infty$ in the unbroken phase, instead. We must therefore look within the more general class of scalar-tensor theories, with actions of the form:

$$S = \int d^4x \, \sqrt{-g} \left(\frac{f(|\phi|)}{16\pi G_0} R - \frac{1}{2} (\partial_\mu \phi)(\partial^\mu \phi) - V(\phi) \right) . \tag{8}$$

This is the action in the frame in which there is minimal coupling between matter and gravity, pinning down a physical frame. For this model, we have [24]:

$$G_{eff} = G_0 \frac{1}{f} \frac{f + 4f'^2}{f + 3f'^2} \sim \frac{G_0}{f(|\phi|)} \tag{9}$$

so that

$$G_{eff} \sim \frac{1}{f} \to 0 \tag{10}$$

if $f \to \infty$ as $|\phi| \to 0$. When $\phi = \phi_0$ (the minimum of the 'Mexican hat' potential) we should have $f = 1$. Any such function realizes our idea. However turning off the coupling between matter and gravity is quite different from turning off all of gravity's dynamic degrees of freedom. Gravitons still exist. The complex array of vacuum solutions known to exist in scalar tensor theories, both isotropic and anisotropic, are still possibilities for the early universe.

These models are potentially very interesting for modelling the early universe, and discussing various potential cosmogonic scenarios, as we now illustrate. First, consider a crude cosmological model in general relativity (GR) with an interval of time during which $G = 0$ for $t_1 < t < t_2$ and $G = G_0$ otherwise. This simple square-well gravity can be modelled by multiplying the usual fluid energy-momentum tensor in Einstein's equations by the sum of Heaviside functions

$$Y_{12} = 1 - H(t - t_1) + H(t - t_2), \tag{11}$$

where $H(x) \geq 0$ for $x \geq 0$ so $Y_{12} = 0$ for $t_1 < t < t_2$ and zero otherwise. The Friedmann equation is

$$3\left(\frac{\dot{a}}{a}\right)^2 = 8\pi G_0 \rho Y_{12} - \frac{3k}{a^2} + \Lambda. \tag{12}$$

For radiation (with $\rho \propto a^{-4}$) we have $a^2(t) \propto Y_{12}t - kt^2 = -kt^2$ when $\Lambda = 0$ for $t_1 < t < t_2$ and so we must have $k = -1$ for a solution to exist, and then $a(t) = t$ and there is a coasting period of Milne expansion with no particle horizon present. When $\Lambda \neq 0$ we have de Sitter solutions for all k so long as $a^2 > 3/\Lambda$ with $a(t) \propto \cosh(t\sqrt{\Lambda/3})$ for $k = +1, a(t) \propto \exp(t\sqrt{\Lambda/3})$ for $k = 0$, and $a(t) \propto \sinh(t\sqrt{\Lambda/3})$ for $k = -1$. However, during the interval (t_1, t_2) there can still be solutions even when $\Lambda = 0$ and $k \geq 0$ because in reality there will always be departures from exact isotropy which contribute (possibly very small) shear terms on the RHS of eq.(12) proportional to a^{-6} which drive the solution towards $a \propto t^{1/3}$ when $k = 0$, for example.

Again, we see that turning-off G still leaves some gravitational degrees of freedom: it reduces the free data required to specify a general GR cosmological solution on a spacelike hypersurface from 8 arbitrary functions to 4 for a perfect fluid, and from 6 to 4 for a self-interacting scalar field [25]. If there is inhomogeneity, with $t_1(\vec{x})$ and $t_2(\vec{x})$, then $Y_{12}(\vec{x})$ and the duration of this finite period of zero-G evolution, will be spatially varying and create residual density variations at $t_2(\vec{x})$ even if ρ were uniform before t_1.

This type of "on-off-on" gravity scenario could be implemented in the action (8) in a theory with an scalar potential of the form $f(|\phi|) \propto G_{eff} \propto Y_{12}^{-1}$. In this case, we would not simply recover the vacuum solutions on the interval (t_1, t_2), as there are many vacuum solutions in generalized scalar-tensor theories like Eq. (8) with non-constant ϕ, just as there are in the simplest case of Brans-Dicke gravity. The square-well form (11) can now be smoothed to the form $\tanh^{2n}(\lambda\phi)$, with $n \in \mathbb{Z}^+$, to mimic the turn-off effect in (t_1, t_2) during the very early universe. We will examine the interesting consequences of such scenarios elsewhere, but the general picture is one we have already sketched: one in which the early universe is simpler during a finite interval of time, due to the partial switching off of gravity.

4. Conclusions

To conclude, we thus see that a number of scenarios may be contemplated for both the early universe and the puzzle of quantum gravity, depending on how we implement our broad idea that the Higgs mechanism is responsible for switching on gravity. In whatever realization, the Big Bang universe's early stages (and possibly the Big Bang singularity itself) are replaced by a different type of solution. In the first realization proposed here a topological solution would model the early universe; in the second a vacuum solution of a scalar-tensor theory would rule it. In both cases the cosmogony is dramatically modified, with opportunities to remove the cosmological singularity. The possibility that a similar phenomenon might occur in the future would also modify the eschatology, rendering the future of the Universe simpler. We stress that the phenomenon envisaged here is different from that proposed in [30, 31, 32], where the muting of gravity applies only to the cosmological perturbations. However the two ideas could be complementary.

One further scenario we wish to pursue in the future, is based on the conjecture that, as a converse to the expectation that all symmetries are restored at high enough energies, perhaps they are all broken at low enough energies. This has been discussed in connection with the breaking of $U(1)$ symmetry at low energy, thus creating a photon mass and even the possibility of photon oscillations via $U(1) \times U'(1)$ (see [26, 29]). If such symmetry breaking in the future drove $G \to 0$, then t_1 might exceed the present age and $t_2 = \infty$. The future cosmological evolution of unbound structures at $t > t_1$ would not be significantly altered, since it is already destined to be dominated by G-independent factors, like curvature, frozen-in anisotropies with constant shear to Hubble ratio, or Λ, in the future of open universes (cf. 12)). However, all bound structures would unbind and black hole horizons would shrink to zero. In this sense the future would be asymptotically simpler.

Regarding the conundrum of quantum gravity, we recall that most of the work attempting to solve the puzzle has been performed in the absence of matter fields. It has been speculated that the problems with the UV limit of the theory could be resolved with the addition of matter. This could indeed be the case in the scenarios considered here: specifically the Higgs field could turn off gravity at high energies, either trivializing its UV limit by rendering it topological (as in [33]), or switching off its interactions and self-interactions. This would be the ultimate asymptotic freedom. The full implications of this scenario should be studied in more detail, but it should be clear that this is a new avenue of enquiry.

We close with two interesting avenues of further study for this type of models. As with any phase transition with spontaneous symmetry breaking we know that a cosmic network of topological defects may be produced, with a morphology dependent on the homotopy groups of the quotient of the full group and the broken subgroup. For example, for a real scalar Higgs field one would have domain walls. Inside the topological defect the symmetry remains unbroken, even after the transition, the field remaining stuck on the false vacuum. The implication, in our scenario, is that gravity would remain switched off at the core of any defect associated the symmetry breaking responsible for turning on gravity. Such solutions could be very interesting, both theoretically and phenomenologically. On a different front the weak field limit of these theories is also bound to contain interesting phenomenology. This was already discussed in the past in [4]. Corrections to the Newtonian potential will be more involved than just a shift in G or of the Yukawa type. We defer to a future publication the full investigation of the ensuing constraints.

Acknowledgements

JDB is supported by STFC and also acknowledges support and hospitality from the Big Questions Institute, the School of Physics, and Faculty of Science at the University of New South Wales, Sydney during this work. JM acknowledges support from John Templeton Foundation, a STFC consolidated grant and the Leverhulme Trust.

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