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# Symmetry Breaking and Newtonian Gravity

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Abstract: We study electron-electron interaction via resonant Higgs exchange in the presence of electroweak-vacuum muon loops. We treat the electroweak Higgs boson in a fully quantum-mechanical way, and obtain a "microscopic" basis for the static Newtonian force law of gravity. We obtain a simple relation between Newton's constant and the mass of the Higgs boson, which we estimate to be 85 GeV (with a rather narrow decay width to $\bar{b}b$ of about 4 MeV).

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In this paper we wish to study the possibility of determining, in an entirely heuristic manner, a quantum field-theoretic origin for the gravitational interaction [1-3]. We will ascribe the term "microscopic" to this endeavour; however, we will see that this term has an imprecise meaning in the present context, just as it does in the description of the Casimir effect [4], in which "microscopic" vacuum fluctuations are inextricably influenced by macroscopic boundary conditions. In order to isolate the essential elements of the problem we have set ourselves, we find it convenient to quote a paragraph from Ref. [5], pp. 6-7: "The symmetry" of an interacting quantum field theory "under the group of inhomogeneous Lorentz transformations is usually assumed to be an exact symmetry, at least so far as elementary particle physics is concerned. If this symmetry is exact, then space is flat. But according to the theory of general relativity, it is possible for space to be curved. Perhaps then, rather than considering the Poincaré group, we should consider a larger group, of which the Poincaré group is only an asymptotic limit. However, we shall not consider this possibility." Instead, we will adopt the view that a field theory with space-time curvature is a "long-range" phenomenological theory that is based on an underlying "microscopic" theory in which space-time is flat but distorted in a topological sense by macroscopic boundary conditions associated with the electroweak Higgs mechanism [6]. By choosing to work in a flat space-time and retaining exact Poincaré invariance, we will easily be able to use the language of quantum field theory in our attempt to initiate a broader understanding of the foundations of the gravitational interaction. We will restrict ourselves to a phenomenological study of the elements required to derive the static Newtonian force law. Our level of rigour will be such that it will be useful to keep in mind a remark attributed (see Ref. [7], pg. 43) to E. Teller in the 1930's: "I don't understand it, but I will explain it to you."

The starting point for our investigation is provided by the observation that the Yukawa coupling between an electron and an electroweak Higgs boson is very small; this creates the possibility that a massless two-Higgs bound state could mediate a long-range force of extremely small strength. We will therefore study electron-electron interaction via two-Higgs exchange, with the aim of determining the minimal phenomenological considerations needed to provide the desired two-Higgs binding. We will make the somewhat unusual assumption that a pair of electroweak-vacuum muon loops that exist in balance with electroweak symmetry breaking are sufficient to "bind" the Higgs particles in the sense of resonant alignment of the underlying non-perturbative vacua. We will use Fig. 1 as a framework for our investigation of a consequence of this assumption; a proof of the assumption itself has to be left to a sufficiently interested reader. We use the Euclidean-space notation of Ref. [8]. We note that it would be interesting to see if some of our considerations might prove to be relevant, in an approximate sense, to the study of pion-pion scattering in low-energy quantum chromodynamics (QCD) [9], where the pion is the relevant Goldstone boson. It is also worth noting that the study of meson decays in three space-time dimensions may be of relevance to condensed-matter physics [10], and might provide a context in which some of our considerations could be applied to the study of resonance-like behaviour associated with $e^+e^-$ peaks observed in heavy-ion collisions [11].

Fig. 1

We begin by assuming that electroweak symmetry breaking has taken place in a largely unspecified manner, with the result that we have a Higgs boson of mass $m_H$ that interacts with an electron of mass $m_e$ via a Yukawa coupling of strength (see Ref. [12], Appendix F)

$$g_{He} \equiv \frac{e}{2 \sin \theta_W} ( m_e/m_W ) , \tag{1}$$

where $\alpha = e^2/(4\pi)$ is the fine-structure constant, $\theta_W$ is the Weinberg angle, and $m_W$ is the mass of the W boson. We further assume that chiral symmetry has been broken in the QCD sector, and we wish to specify an effective Yukawa coupling $g_{Hq}$ between the Higgs boson and a dynamical quark propagator in a hypothetical generalization of Fig. 1. In order to do this, we will use a recent generalization of the Nambu — Jona-Lasinio mechanism of dynamical mass generation [13] to the renormalizable theory of QCD [14-16]. (We have also used this method to study dynamical mass generation in the Gross-Neveu model [17] at finite temperature and density [18].) In this approach, we write the (massless) fermionic part of the QCD Lagrangian, $L_{quark}$, as the sum of a "perturbative" part and a "counter-term" part, i.e.

$$L_{quark} \longrightarrow L_{pert} + L_{ct} , \tag{2}$$

where

$$L_{pert} = L_{quark} + i\, m\, \bar{\psi}\psi$$

and

$$L_{ct} = -i\,(\delta m)\,\bar{\psi}\psi .$$

We evaluate the one-loop fermion mass renormalization $\Sigma_m$ using $L_{pert}$ in (2), and require that the limit $\delta m \longrightarrow m$ be taken so that the original Lagrangian remains massless; the dynamical quark mass will then be a quantum result obtained by solving the "gap equation" $\Sigma_m = 0$ in a manner consistent with the leading-logarithm approximation of renormalization-group-improved perturbation theory (see Ref. [16] for detailed discussion). We obtain an off-shell (virtual) dynamical quark mass $m(p^2)$ specified by [14-16]

$$m\, Z^{1/2} = \exp[2/3]\, \Lambda , \tag{3}$$

where

$$Z = (1 + p^2/m^2)^{1 + m^2/p^2}$$

and $\Lambda$ is a one-loop renormalization-group invariant. At the nominal pole

position $p^2 = -m^2$ , we obtain from (3) the "physical" (real) quark mass $m_0 \equiv \exp[2/3]\Lambda$ ; we note that $m(p^2)$ vanishes for $p^2 \geq m_0^2$ . If we choose $\Lambda = 175$ MeV [16], we obtain $m_0 \approx 340$ MeV; this is the value of the constituent quark mass that follows from additivity of quark magnetic moments inside the nucleon [19]. Previous study [20] of the behaviour of liquid helium-4 might provide a useful analogy (see Ref. [16] for a brief discussion) to help us understand how the physics of strongly interacting quarks and gluons may effectively be reduced to the study of "weakly" interacting constituent quarks moving in a confining potential, although our specific phenomenological approach [14-16] may still need justification. For the time being, we will only make use of the method of derivation of the "physical" quark mass $m_0$ , and we refer the interested reader to Refs. [14,16] for additional discussion and references.

We interpret the rearrangement of the Lagrangian (2) in terms of a balance between a tree-level mass of electroweak origin and a QCD counter-term that represents the response of the QCD vacuum to this tree-level mass (see Ref. [16]); this balance ensures not only that the electroweak mass m does not obtain radiative corrections within QCD, but that its value is dynamically determined in a way that enables us to exchange a free electroweak parameter for one that may be determined experimentally purely within QCD. In particular, we see that our desired effective Higgs-quark Yukawa coupling $g_{Hq}$ is specified by (1) with $m_e$ replaced by $m_0$ . This indicates that we can, in principle, generalize the interaction of Fig. 1 to include nucleons in our flavour-independent approach.

We note that we intend Fig. 1 to represent a physical two-Higgs vacuum in the sense that the effective electron "ends" of the process are to be thought of as quantum-mechanical "hooks" to which the external fermion legs and an

intervening "tree-level" propagator describing the exchange of an effective two-Higgs bound state may be attached, thereby effecting momentum transfer. It is worth noting, in the context of the work of this paper, that the conclusion of Ref. [21] tells us that a hypothetical scalar particle that exists in conjunction with the usual products of beta decay is not observable in terms of a final state that is independent of the underlying structure of the physical vacuum associated with that particle. We use this to conclude that our "tree-level" two-Higgs bound state exists only in the sense of a virtual excitation of the physical vacuum, specified by Fig. 1, on which we superimpose it. As a result, we must also emphasize that there is no such thing as a "single" (two-Higgs bound state)-electron vertex, and that Fig. 1 therefore represents a somewhat complicated phenomenon. In this paper, we will examine a phenomenological consequence of assuming that the effective two-Higgs bound state that we superimpose on the graph of Fig. 1 is massless at "tree-level." We note that the possible existence of such a massless two-Higgs state in the presence of electroweak-vacuum muon loops may be analogous to the existence of Josephson tunneling in superconductivity [22].

In order to study the non-perturbative vacuum associated with the graph of Fig. 1, we must treat the Higgs boson in a fully quantum-mechanical manner. We choose to accomplish this in the present "low-energy" setting by using a simple trick borrowed from study of the Gross-Neveu model in the auxiliary-field formalism [17]. We write the Higgs-fermion part of the Lagrangian, $L_{Hf}$, in the simple flavour-independent form

$$L_{Hf} = \frac{1}{2}\partial_\mu\phi_0\partial_\mu\phi_0 + \frac{1}{2}m_B^2\phi_0^2 + i\,g_0\phi_0\bar{\psi}\psi \ , \tag{4}$$

where $m_B$ is the formally infinite "bare" off-shell contribution to the perturbative Higgs mass arising from interaction of the Higgs with an "explicit"

electroweak vacuum condensate (the off-shell nature of the physical Goldstone boson has been discussed in Ref. [16] for the case of the pion, and we assume a similar interpretation for the physical Higgs particle). Following Ref. [17], we choose to define the "physical" Higgs field and Yukawa coupling by $\phi = m_B \phi_0$ and $g = g_0/m_B$ , respectively (to make sense of the units, we prefer to insert factors of a unit mass where necessary). Taking the limit $m_B \longrightarrow \infty$ , the Lagrangian (4) becomes

$$L_{Hf} = \frac{1}{2} \phi^2 + i g \phi \bar{\psi} \psi \quad . \tag{5}$$

The Lagrangian (5) is a realization of the statement that we may regard the non-perturbative Higgs boson as a quantum mode because Goldstone's theorem [23] is a quantum result. The Higgs kinetic term disappears from the classical Lagrangian, and the physical mass $m_H$ of the Higgs boson is "absorbed" into the underlying non-perturbative "one-Higgs" vacuum, of which the physical Higgs boson is rigorously a "zero-energy" excitation. We assume that resonant anti-alignment of the "one-Higgs" vacua in the presence of the muon loops in Fig. 1 is the mechanism responsible for the masslessness of our "two-Higgs" bound state; this also allows us to ignore the details of the "two-Higgs" vacuum [24] in our "auxiliary-field" approach.

The Lagrangian (5) specifies simple "perturbative" Feynman rules for the one-Higgs lines and couplings in Fig. 1. However, because our physical Higgs boson is rigorously a "zero-energy" excitation, there is no Higgs-muon -Higgs-muon loop integration. In addition, the non-perturbative one-Higgs vacuum is not smoothly matched to the perturbative vacuum represented by the muons in the loops of Fig. 1, so that a vacuum decay process must occur at each of the Higgs- muon vertices. The fermion propagators between the Higgs-fermion vertices therefore correspond not to virtual, off-shell states but to real, off-shell states that are produced and then absorbed in a resonant "external"

exchange between "adjacent" vacua. The muon propagators must absorb the mass-squared, $q^2 = m_H^2$ , associated with the non-perturbative one-Higgs vacuum; in accord with our desire to isolate vacuum properties, we neglect external momenta because we have specified that transfer of external electron momentum occurs only via the "perturbative" two-Higgs propagator that we superimpose on Fig. 1. The ambiguity in the sign of the square root of $m_H^2$ requires us to replace the simple virtual muon propagator by the real, off-shell state specified by the anti-symmetric (for fermion statistics) wave-function

$$-i\,2^{-1/2}\,[\,\bar{u}(p_+)(\not{p}_+ + m_\mu)^{-1}u(p_+) - \bar{v}(p_-)(\not{p}_- + m_\mu)^{-1}v(p_-)\,]\ , \tag{6}$$

where $p_+^2 = p_-^2 = m_H^2$ , the momenta $p_+$ and $p_-$ refer of course to positive and negative energy states, respectively, and the spinors u and v satisfy the (real, off-shell) relations (see Ref. [25], pp. 75-80)

$$\bar{u}(p_+)\,u(p_+) = -\bar{v}(p_-)\,v(p_-) = 2\,m_H$$

as well as the relations

$$\bar{u}(p)\,\gamma_\mu\,u(p) = \bar{v}(p)\,\gamma_\mu\,v(p) = 2\,p_\mu\ ;$$

the factor of (-i) in (6) comes from the propagators sandwiched between the spinors.

The expression (6) reduces to the effective "vertex" factor specified by

$$\begin{aligned} f(m_\mu, m_H) &\equiv -i\,2^{5/2}\,m_\mu m_H/(m_H^2 + m_\mu^2) \\ &\approx -i\,2^{5/2}\,(m_\mu/m_H)\ , \end{aligned} \tag{7}$$

where we have included a factor of 2 for the two "directions" in which the Higgs mass-squared can be sent through the perturbative muon line that it joins. After we superimpose a $1/q^2$ propagator for the "tree-level" exchange of a massless two-Higgs bound state, and include a combinatoric factor of 36 for the ways in which the fermion lines may be joined by the Higgs lines, the

amplitude specified by the graph of Fig. 1 becomes

$$A(q^2) = (36\, q^{-2})(ig_{He})^2(ig_{H\mu})^6[\, f(m_\mu, m_H)\,]^6$$
$$\equiv -16\pi G_0\, m_e^2\, q^{-2}\ , \tag{8}$$

where $g_{H\mu}$ is the muon-Higgs Yukawa coupling given by replacing $m_e$ with $m_\mu$ in (1). (We note that, in our units [8], the Feynman-gauge amplitude for photon exchange, $4e^2q^{-2}$, yields the Coulomb potential $\alpha/r$.)

If we tentatively identify $G_0$ in (8) with Newton's constant $G_N$ given by (see Ref. [26])

$$G_N = 6.70711(86) \times 10^{-39}\ \mathrm{GeV}^{-2}\ ,$$

we obtain a simple relation between $G_N$ and $m_H$, namely

$$G_N \approx \frac{288}{\pi}\,[\,4\pi\alpha/\sin^2\theta_W\,]^4\, m_\mu^{12}\, m_W^{-8}\, m_H^{-6}\ ;$$

use of the standard values [26] $\alpha \approx 1/137$, $\sin^2\theta_W \approx 0.233$, $m_\mu \approx 0.106$ GeV and $m_W \approx 80.2$ GeV then yields the estimate

$$m_H \approx 85.3\ \mathrm{GeV}\ . \tag{9}$$

The estimate (9) is well within the range of present expectations for the Higgs mass [27], and may be easily susceptible to experimental verification.

It would be interesting to see if our approach to the study of the static Newtonian gravitational interaction might be of relevance to the study of CP violation [28], which has been tentatively linked [29] to the possible existence of "anti-gravity," i.e. opposite sign for matter-antimatter gravitational interaction as for matter-matter interaction. We note as well that the work of Ref. [30] might provide hints as to how flavour-dependent effects might be incorporated into the dynamical approach of Refs. [14-16]. The gravitational effect that we have studied will be affected by the observation that the masses of constituent quarks in nucleons will vanish by means of a second-order chiral-symmetry-restoring phase transition at a

critical temperature of about 200 MeV [15]. We wish to keep in mind, however, that the pure electroweak (i.e. flavour-dependent) "current" quark masses that do not enter into the flavour-independent analyses of Refs. [14-16], and are unrelated to the phenomenon of quark confinement (cf. Ref. [30]), will not be significantly affected by this phase transition (we feel that this conclusion is implied by the work of Ref. [31] because of their "perturbative" treatment of the $\delta m$ term, see Ref. [15] for discussion). In any event, the effects of finite temperature on the gravitational interaction are worth studying in detail (see Ref. [32]).

Because the study of CP violation is important in its own right, we wish to advocate consideration of any viable projects that can carry out such study in detail, including the possibility of a high-intensity linac intersecting with a storage ring, as well as the study of proton-(anti)proton scattering at sufficiently high luminosity (see also Ref. [33]). It is of interest in regard to study of the b quark that the result (9) tells us that the dominant Higgs decay will be $H \longrightarrow \bar{b}\, b$ , with a narrow partial width given by (see Chanowitz [24], pg. 387, and Thorndike [33], pg. 227) $\Gamma(H \rightarrow \bar{b}\, b) \approx (1.6 \times 10^{-4}\, \mathrm{GeV}^{-1}) m_b^2$ $\approx (3.8 \pm 0.3)$ MeV. Finally, we note that the one-Higgs vacuum decays in Fig. 1 may also be reminiscent of the Drell-Yan process [34] in hadronic scattering; it would be interesting if this possible analogy were to help in visualizing the detailed processes occurring in Fig. 1. It would be particularly interesting to consider the origins of corrections to the static Newtonian force law.

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We point out, perhaps belatedly, that our implicit use of the word "scalar" in the context of "tree-level" two-Higgs exchange superimposed on the graph of Fig. 1 should be interpreted in an "off-shell" sense, corresponding to use of the harmonic gauge in a linearized treatment of quantum gravity mediated by exchange of an effective spin-2 disturbance (the "graviton"; see Ref. [25], pp. 282-4, especially Eqn. (14.19)). It is possible that an appropriate quantum-mechanical description of light bending (cf. Ref. [25], pp. 291-3) could be obtained in analogy with the photo-electric effect [35] by having an electron in the electroweak vacuum kicked or spin-flipped by the photon, so that the resulting "displacement" with respect to its vacuum surroundings would expose the resulting off-shell electron state to direct interaction with an "auxiliary-field" Higgs (cf. Fig. 1) before the electron re-emits the photon and "disappears" back into the electroweak vacuum. We leave these, and other, problems of interpretation, analysis and of course verification, to the interested reader.

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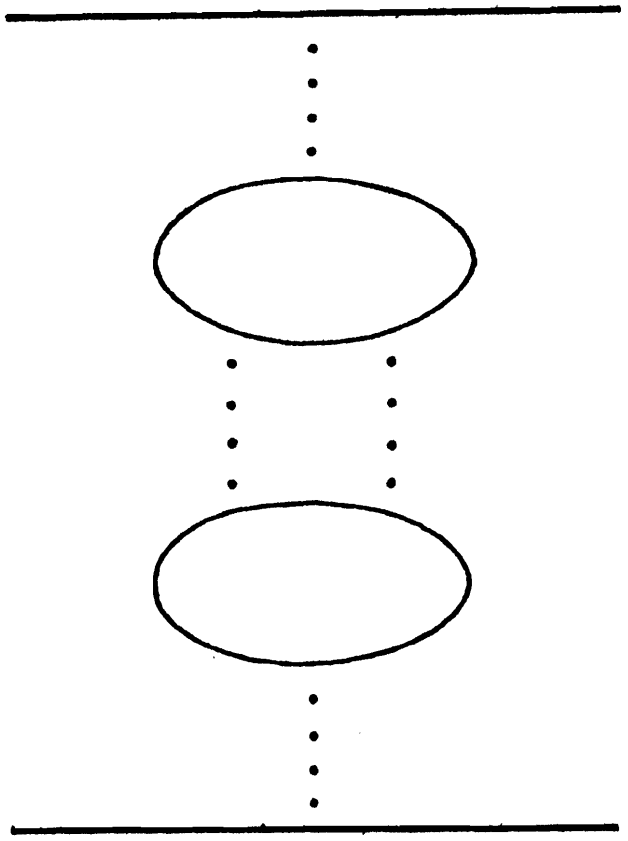

Fig. 1. This diagram allows us to study electron-electron interaction via resonant Higgs exchange in the presence of electroweak-vacuum muon loops.