A New Ether Born out of Quantum Gravity?

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Reconsidering one's position on preferential rest frames (P.R.F.)
Emphasize: It is very hard to write a theory that violates Lorentz invariance (L.I.) which does not require the notion of a preferred rest frame (P.R.F.)

This does not mean one can not conceive of such things: [in a context as QG imagine a certain "vacuum state" that selects a P.R.F. but then consider the state of the world to be a quantum superposition of such states for all "frames".]

The point is nonetheless that it is very difficult to think (or write) how such breakdown of L.I. might become manifest!

So in dealing with Q.G. phenomenology one is in practice always invoking the existence of a P.R.F. →

"The resurrection of the Ether"?

Nowadays the P.R.F. is identified with the one selected locally by the C.M.B.

Could we think about the "Dragging of Frames"?
The earth moves w.r.t. that P.R.F at about 300 m/s, so we can convert modern probes of such motion into tests of Q.G. induced breakdown of h. I of the type that is being considered nowadays. One of the first considerations (within h. Q.G) was introduced by Gambini & Pullin who studied the Hamiltonian constraint in the presence of matter.

\[ E^{\alpha}\partial_{\alpha}E^{\mu}F_{\mu\nu}(A)_{\nu} + \hbar(A^{\alpha}_{\beta}E^{\beta}_{\alpha}; \Phi) = 0 \]

geometric variables \quad matter fields.

From this, the matter Hamiltonian is identified as:

\[ \hat{H} = \int h(A, E, \Phi) \, d^3x \]

Then, upon selection of a fixed appropriate state for the geometrical D.O.F: A, E, the effective Hamiltonian for the matter is taken to be

\[ \hat{H} = \langle \text{weave} | \hat{H} | \text{weave} \rangle \]

for matter = Maxwell field G.P. obtain

\[ \hat{H} = \int d^3x \, E^2 + B^2 + \frac{\hbar}{M_p} (E \partial_x E + B \partial_x B) \]

which breaks L.I. Purity ... It describes exotic 'time' of P.R.F.
The corresponding Lagrangian can be written as

\[
L = \mathcal{F}_{\mu\nu} F^{\mu\nu} + \frac{5}{M_p} W^m F_{\mu\nu} W^m \left( \omega^m F^{\mu\nu} \right)
\]

where \( W^m \) is the four velocity of the P.R.F.

Similarly, following Alfaro, Madsen, Unruh, one writes for the fermions

\[
\mathcal{L}_{\text{fermions}} = \frac{1}{2} \bar{\psi} \left( i \gamma \gamma - m \right) \psi + \frac{\Theta_1}{M_p} \bar{\psi} \gamma^m \left( \gamma^m - i \omega^m \right) \psi
\]

\[
+ \frac{\Theta_2}{M_p} \bar{\psi} \gamma^m \gamma^n \gamma^o \gamma^p \psi \gamma^m \frac{\partial}{\partial \gamma^m} \psi - \frac{1}{4} \Theta_4 m^2 \bar{\psi} \gamma^m \gamma^n \gamma^o \gamma^p \psi \gamma^m \frac{\partial}{\partial \gamma^m} \psi
\]

Interactions can be introduced by minimal coupling. The point is that in the lab \( W^m \) has spatial component that changes direction as a result of the Earth's rotation. In fact the effective 2-component (low energy) Hamiltonian for a fermion in the lab becomes

\[
H_{\text{eff}} = m \left[ 1 + \Theta_1 \frac{m}{M_p} (\vec{W}^2) \right] + \left[ 1 + 2 \Theta_2 \frac{m}{M_p} \left( 1 + \frac{5}{3} (\vec{W}^2) \right) \right] \left[ \frac{\vec{p}^2}{2m} + g \mu \vec{S} \cdot \vec{B} \right] + \left( \Theta_2 + \frac{1}{2} \Theta_4 \right) \frac{m}{M_p} \left[ \left( 2m - \frac{2 \vec{p}^2}{3m} \right) \vec{S} \cdot \vec{B} \right]
\]

\[
+ \frac{1}{m} \vec{S} \cdot \vec{Q} \cdot \vec{W} \right] + \Theta_1 \frac{m}{M_p} \left[ \frac{\vec{W} \cdot \vec{Q} \cdot \vec{W}}{m} \right]
\]
The point is that modern experiments set strong bounds on such couplings (tests of isotropy of physical laws).

**Basic Idea:** Consider a $J = \frac{3}{2}$ nucleus in an external $B$

![Diagram](image)

$$M = \begin{cases} 
\frac{3}{2} & \\
\frac{1}{2} & \\
-\frac{1}{2} & \\
-\frac{3}{2} & 
\end{cases}$$

Isotropy $\Rightarrow$ equal splitting of lines

$\delta$ would lead to unequal splitting which oscillates during the day!

Fourier transform of data removes noise and lets one focus on the desired effect, leading to very precise bounds.

This leads to [Chupp et al PRL 63 (1989) 1541, 12 hr period]

$$|\theta_1| < 3 \times 10^{-5}$$

[Similar to astrophysical bound obtained by Gleiser & Kosmich.]

and [Phillips et al PRD 63 R111101 (2001), 24 hr, protons]

[Bean et al PRL 85, 5038 (2001), 24 hr, neutrons]

$$|\theta_2 + \frac{1}{2} \theta_4| < 2 \times 10^{-9}$$

[Similar to Jacobson et al but constraining both signs!]

(Synchrotron radiation from He)

(CEAB module)
We can in fact do better by noting that although very high energy particles are not so easily found (high energy cosmic rays & photons from astrophysical sources ...), virtual photons (and other particles) with arbitrarily high energy contribute to ordinary processes. QG modifications would also affect these virtual particles (in a sense all particles are virtual...). We can consider then the effects of L-V in the propagation of a virtual photon, on an electron's self energy

\[ \text{Modify photon propagator.} \]

\[ = \text{standard self energy} \]

\[ + e^2 \frac{\bar{\psi} \gamma^\mu \psi}{M_p^2} \frac{1}{(p - k)^2 - m^2} \frac{1}{M_p^2} \]

diverges quadratically, but of course we should regard our theory as effective & valid only below \( M_p \).

This leads to a term

\[ \frac{e^2 \frac{\Lambda^2}{M_p^4}}{M_p^2} \bar{\psi}(p) i \sigma^\mu \nabla^\mu \psi(p) \]

which is a radiative correction to \( \Theta_4 \) so

\[ \frac{\Theta_4}{M_p^2} = \frac{\Theta_4}{M_p^2} + \frac{3}{m^2} \frac{e^2 \Lambda^2}{M_p^2} \]

The factors \( 1/M_p^2 \) cancel out!

But \( \Lambda^2 \sim M_p^2 \) remains!
In the case of the photon's self-energy, the effects are less dramatic:
\[ \frac{1}{2} \Phi^{\mu \nu} = 3 B + e^2 \sum_i c_i \ln(N_i) \]
Setting for instance \( \Lambda = M_p \left( \frac{10^{-2}}{\ln(10^{20})} \right) \approx 40 \)
But our bounds for the fermions become
\[ |\phi| < 10^{-40} \]
Myers & Pospelov have noticed this problem which would be very serious even if one sets \( \Lambda = \Lambda_{\text{susy}} \approx 10^{-10} \text{GeV} \)
It is not clear however what would this mean? ?

They have proposed a recipe (ad hoc) that they suggest could avoid this fatal problem: The idea is implemented within the context in which, at the effective F.T. level the Q.G. effects are described by Dimension 5 operators (non-renormalizable).
So they write:

\[ L_{\text{maxwell}} = \frac{3}{M_p} W^\mu F_{\mu\nu} W^\nu \phi (W^\nu \phi) \]

\[ L_{\text{scalar}} = \frac{k}{M_p} \phi (W^\mu \phi)^3 \phi \]

\[ L_{\text{fermion}} = \frac{1}{M_p} \left( \bar{\psi} \gamma_{\mu} \left( \gamma^\mu \gamma^5 \psi \right) \right) (W^\mu \phi)^2 \phi \]

And the recipe to avoid the problems is to replace everywhere the tensor \( W^\mu W^\nu W^\rho \) by

\[ C^{\mu\nu\rho} = \frac{1}{6} (\tilde{\psi} \gamma^\mu \psi + \tilde{\psi} \gamma^\nu \psi + \tilde{\psi} \gamma^\rho \psi) \]

which has the property that

\[ C^{\mu\nu\rho} \rho_{\mu\nu} = 0 \quad (\text{for every pair of indices}) \]

Integrals such as

\[ \int \frac{k^\mu k^\nu}{k^4} \text{d}^4 k = \delta^\mu^\nu \Lambda^2 \]

would give zero when contracted with \( C^{\mu\nu\rho} \).

At first sight this seems to work, but it would be very surprising if such a simple recipe were to violate the generic expectations from Effective Field Theory: that all terms compatible with the remaining symmetry would be generated, and thus for consistency one should introduce them from the start.
Now renormalizable terms would be suppressed by appropriate powers of $\Lambda$, but renormalizable terms would not... and there is no reason for them to be absent.

We looked at the issue in some detail and found that higher order (in $\frac{1}{\Lambda}$) terms in the 1 loop calculation of the fermion self-energy (Yang-Mills Theory)

\[
\frac{\text{P-LH}}{P} \Rightarrow \frac{e^2}{(2\pi)^1} \int d^4k \frac{k^2 + m^2 + \frac{2}{M^2} C_{\mu\nu} k^\mu k^\nu}{k^2 - m^2 + i\epsilon} [1 - \frac{1}{k^2 - m^2 + i\epsilon} L - \frac{C_{\mu\nu}}{k^2 - m^2 + i\epsilon} L]
\]

\[
= \frac{e^2}{(2\pi)^1} \frac{2 \frac{\mu^2}{M^2} \int d^4k \frac{l^2 k^\mu k^\nu}{k^2 - m^2 + i\epsilon} C_{\mu\nu} C_{\mu\nu} C_{\gamma\delta\epsilon}}{M^3} + ...
\]

\[
= \frac{e^2}{(2\pi)^1} \frac{2 \frac{\mu^2}{M^2} \int d^4k \frac{l^2 k^\mu k^\nu}{k^2 - m^2 + i\epsilon} C_{\mu\nu} C_{\mu\nu} C_{\gamma\delta\epsilon}}{M^3} + ...
\]

\[
= \frac{e^2}{(2\pi)^1} \frac{2 \frac{\mu^2}{M^2} \int d^4k \frac{l^2 k^\mu k^\nu}{k^2 - m^2 + i\epsilon} C_{\mu\nu} C_{\mu\nu} C_{\gamma\delta\epsilon}}{M^3} + ...
\]

Leading to a $h^2 \mu^2 L^4$ enhanced effective term (supernormalizable)
It is noteworthy that the sort of effects obtained in QG --- as well as the humble String Theory Approach --- inspired models, can be found by a naïve approach based on Q.F.T.

If we assume a Q.G induced granularity of spacetime associated with a P.R.F. --- the rest frame of the D-particle gas ... or the frame where the fluctuations of geometry take a particularly simple form ... (if such frame exists) --- then it is natural to impose a momentum cut-off adapted to that P.R.F.

The idea is to consider loop corrections but keep the virtual's particle momentum components bounded.

Example: Yukawa Theory. If a fermion of mass \( m \), \( \phi \) a massless scalar, and the interaction \( e \bar{\psi} \psi \phi \). Carry out all calculations in the P.R.F. and bound the 3-momenta of the virtual particles \( |\vec{p}| < \Lambda \approx M_p \).

If it is indirectly bound ... binding it directly leads to more severe problems ... unitarity.
We do not expect alternative ways to set the bounds to change things much regarding L.I. Violations because there are no compact regions in four momentum space that are invariant under the Lorentz group.

For diagrams with 2-virtual particles with momenta $\mathbf{h}^n \& \mathbf{P}^m - \mathbf{h}^m$ the integration $d^3h$ becomes

$$\int d^3h \rightarrow 2\pi \left\{ \int_{-1}^{1} d^2h \int dt \left[ \frac{1}{\sqrt{l^2 + \mathbf{P}^2}} \right]^1_{-1} \right\}$$

where $t = \cos (\theta_{\mathbf{h}, \mathbf{P}})$.

Then, the fermion self-energy is

$$\rightarrow_{\text{usual L.I terms}} + \frac{e^2}{64\pi^2\Lambda} \sqrt{(\mathbf{w}_1 - \mathbf{w}_2)(\mathbf{p}_1 + \mathbf{p}_2)}$$

while the boson self energy is

$$\rightarrow_{\text{usual L.I terms}} - \frac{ie^2\Lambda}{64\pi^2} \sqrt{(\mathbf{w}_1 - \mathbf{w}_2)(\mathbf{p}_1 + \mathbf{p}_2)}$$

So, if $\Lambda$ is large, the L.I. in the fermion is small but that in the boson is large if vice versa.
Considering the more relevant QED case, the previous recipe leads to problems with gauge invariance. One variation that has not led to obvious problems in that regard is to cut only the photon's momenta. In this way, there is no connection at loop in the photon's vacuum polarizability and in the electron self-energy we find

\[ e \sqrt{\frac{2 \lambda}{(2\pi)^4}} \frac{\delta^{\mu} P + m}{(2\lambda)^2} \gamma^\nu \frac{\Im \nu}{\lambda^2 + i\epsilon} \]

\[ |\vec{p}| < \lambda \]
\[ \lambda \in (-\infty, \infty) \]

= usual $e$ terms \[ -\frac{i e^2}{24 \pi^2} (\mathcal{M}_\mu - W_\mu W_\nu) \delta^{\mu \nu} \]

[divergent as $\ln \frac{1}{m}$] \[ \uparrow \text{Not suppressed!} \]

Leading to a hard $h \cdot V$ term

\[ \frac{i e^2}{24 \pi^2} \mathcal{F} (\mathcal{M}_\mu - W_\mu W_\nu) \gamma^\nu \mathcal{F} \]

of the form as the Parity Conserving one found by...

Alfonso Woods, Cavenda
Conclusion

Not only are the bounds on possible L I Violations exceedingly tight -- that even a suppression by $E_{\text{mp}}$ can not naturally explained -- but it has become also very hard to understand how would these effects remain so suppressed even if initially one sets them to be so.

In our opinion the spirit of Michelson Morley's result has returned with increased strength in its "quantum reincarnation" and as a result the specter of the Ether—having been summoned using QG inspired enchantments—is once again on the run.

[i.e. already considering $\frac{1}{\lambda_{\text{mp}}}$]