On Massive Gauge Theories Beyond Perturbation Theory

Pilar Hernández

Universidad de Valencia and IFIC

In collaboration with Michele Della Morte, arXiv: 1309.3326

Introduction

The Higgs has been discovered and the EWSB sector of the SM seems accurately described by the simplest Higgs mechanism

Englert, Brout; Higgs; Guralnik, Hagen, Kibble; Weinberg

$$V(\Phi) = -\frac{m_H^2}{2}(\Phi^\dagger\Phi) + \frac{\lambda}{4}(\Phi^\dagger\Phi)^2$$

$$\langle\Phi\rangle = \begin{pmatrix} 0 \\ \frac{v}{\sqrt{2}} \end{pmatrix} \rightarrow m_W = \frac{1}{2}gv, \quad m_f = \lambda_f\frac{v}{\sqrt{2}}, \; m_H = v\sqrt{\frac{\lambda}{2}}$$

BROKEN SYMMETRIES AND THE MASSES OF GAUGE BOSONS

Peter W. Higgs
Tait Institute of Mathematical Physics, University of Edinburgh, Edinburgh, Scotland
(Received 31 August 1964)

In a recent note[1] it was shown that the Goldstone theorem,[2] that Lorentz-covariant field theories in which spontaneous breakdown of symmetry under an internal Lie group occurs contain zero-mass particles, fails if and only if the conserved currents associated with the internal group are coupled to gauge fields. The purpose of the present note is to report that, as a consequence of this coupling, the spin-one quanta of some of the gauge fields acquire mass; the longitudinal degrees of freedom of these particles (which would be absent if their mass were zero) go over into the Goldstone bosons when the coupling tends to zero. This phenomenon is just the relativistic analog of the plasmon phenomenon to which Anderson[3] has drawn attention: that the scalar zero-mass excitations of a superconducting neutral Fermi gas become longitudinal plasmon modes of finite mass when the gas is charged.

The simplest theory which exhibits this behavior is a gauge-invariant version of a model used by Goldstone[2] himself: Two real[4] scalar fields φ_1, φ_2 and a real vector field A_μ interact through the Lagrangian density

$$L = -\tfrac{1}{2}(\nabla\varphi_1)^2 - \tfrac{1}{2}(\nabla\varphi_2)^2 - V(\varphi_1^2 + \varphi_2^2) - \tfrac{1}{4}F_{\mu\nu}F^{\mu\nu}, \qquad (1)$$

where

$$\nabla_\mu\varphi_1 = \partial_\mu\varphi_1 - eA_\mu\varphi_2,$$

$$\nabla_\mu\varphi_2 = \partial_\mu\varphi_2 + eA_\mu\varphi_1,$$

$$F_{\mu\nu} = \partial_\mu A_\nu - \partial_\nu A_\mu,$$

e is a dimensionless coupling constant, and the

about the "vacuum" solution $\varphi_1(x) = 0$, $\varphi_2(x) = \varphi_0$:

$$\partial^\mu\{\partial_\mu(\Delta\varphi_1) - e\varphi_0 A_\mu\} = 0, \qquad (2a)$$

$$\{\partial^2 - 4\varphi_0^2 V''(\varphi_0^2)\}(\Delta\varphi_2) = 0, \qquad (2b)$$

$$\partial_\nu F^{\mu\nu} = e\varphi_0\{\partial^\mu(\Delta\varphi_1) - e\varphi_0 A_\mu\}. \qquad (2c)$$

Equation (2b) describes waves whose quanta have (bare) mass $2\varphi_0\{V''(\varphi_0^2)\}^{1/2}$; Eqs. (2a) and (2c) may be transformed, by the introduction of new variables

$$B_\mu = A_\mu - (e\varphi_0)^{-1}\partial_\mu(\Delta\varphi_1),$$
$$G_{\mu\nu} = \partial_\mu B_\nu - \partial_\nu B_\mu = F_{\mu\nu}, \qquad (3)$$

into the form

$$\partial_\mu B^\mu = 0, \quad \partial_\nu G^{\mu\nu} + e^2\varphi_0^2 B^\mu = 0. \qquad (4)$$

Equation (4) describes vector waves whose quanta have (bare) mass $e\varphi_0$. In the absence of the gauge field coupling ($e = 0$) the situation is quite different: Equations (2a) and (2c) describe zero-mass scalar and vector bosons, respectively. In passing, we note that the right-hand side of (2c) is just the linear approximation to the conserved current: It is linear in the vector potential, gauge invariance being maintained by the presence of the gradient term.[5]

When one considers theoretical models in which spontaneous breakdown of symmetry under a semisimple group occurs, one encounters a variety of possible situations corresponding to the various distinct irreducible representations to which the scalar fields may belong; the gauge field always belongs to the adjoint representation.[6] The model of the most immediate interest is that in which the scalar fields form an

The big questions for the last 50 years....

- Any new physics scale beyond the SM, Λ, leads to a **hierarchy problem**

$$m_H^2 \sim \Lambda^2 \log(\Lambda) \quad \text{or} \quad v^2 \sim \Lambda^2 \log(\Lambda)$$

  The EW scale needs to be fine-tunned...

- Even if we declare there is no other scale of physics (at least within the realm of QFT) we face a **triviality problem**

$$\lambda = 0 \rightarrow \frac{m_H}{v} = 0$$

  Seems to indicate that the theory needs new physics...

- Couplings to the Higgs create havoc in the SM: **the flavour problem** ...

In spite of all these problems SM is (compared to its contenders) beautifully simple and predictive in perturbation theory...can it be even more beyond ?

Beyond perturbation theory

This naive characterisation of spontaneous symmetry breaking is not at work beyond perturbation theory

In any lattice formulation of gauge theories: Elitzur's theorem

$$\langle \Phi \rangle = 0$$

Elitzur; Frölich et al

Even if the gauge is fixed, the vanishing of this expectation value depends on the gauge choice

A non-perturbative derivation of the weakly-coupled electroweak sector of the SM might shed some light on the hierarchy, triviality and flavour problems...

Lattice Field Theory

Discretization of a QFT on a euclidean space-time lattice:

$$\partial_\mu \Phi(x) \rightarrow \Phi(x + a\hat{\mu}) - \Phi(x)$$

Our observables are euclidean correlation functions (KL)

$$\lim_{x_0 \rightarrow +\infty} \langle \Phi(x)\Phi(0) \rangle = \sum_\alpha \int \frac{d^4 p}{(2\pi)^4} \frac{Z_\alpha e^{ipx}}{p^2 + m_\alpha^2} \simeq \sum_\alpha \int_{\mathbf{p}} \frac{Z_\alpha e^{-E_\alpha(\mathbf{p}) x_0} e^{i\mathbf{p}\mathbf{x}}}{2E_\alpha(\mathbf{p})}$$

There are infinite ways of discretising a field theory: the physics is uninteresting non-universal physics, unless scaling (critical) regions exist

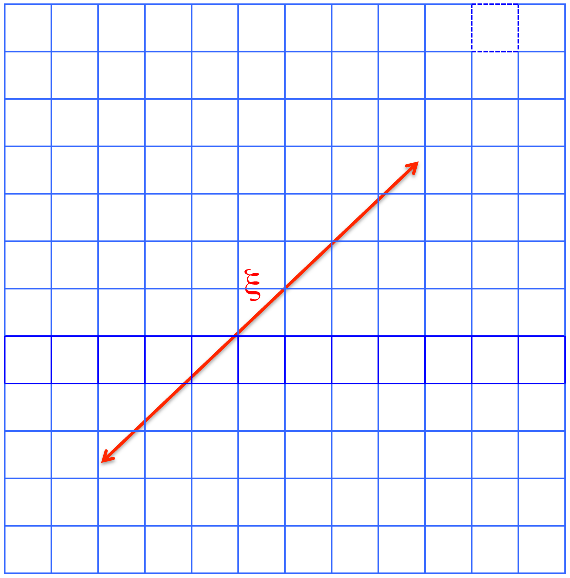
$$(\xi_\alpha / a)^{-1} \equiv m_\alpha a \ll 1$$

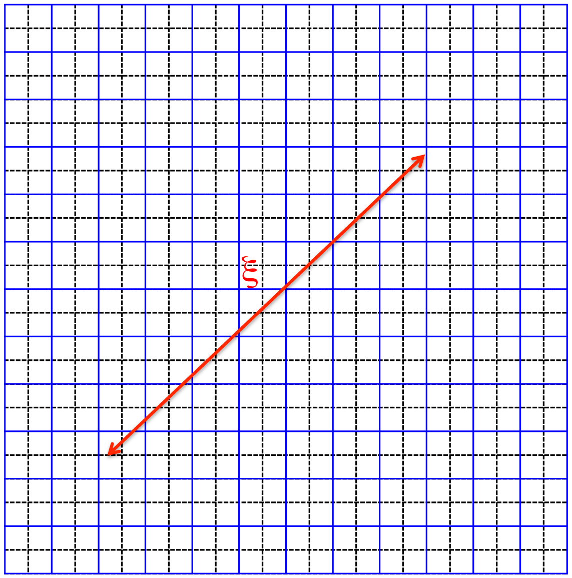
There are very light dofs compared to the cutoff a^{-1}: the dynamics of these fields is universal, described by an effective QFT

Unitarity can be warrantied by the property of reflection positivity

Osterwalder, Seiler

Wilsonian Renormalizability:





$$S^{(n)}(a) = \sum_{i,x} g_i^{(n)} O_i(x) \qquad S^{(n+1)}(a) = \sum_{i,x} g_i^{(n+1)} O_i(x)$$

$$g_i^{(n+1)} = R_i^{(n)}(g^{(n)})$$

Fixed-points of RG: $g_i^* = R_i(g^*)$

$$m_\alpha(g^*) = \text{fixed} \rightarrow m_\alpha(g*)a \rightarrow 0$$

Wilsonian Renormalizability:

Critical Regions $\leftrightarrow$ Effective QFT

Fixed Points $\leftrightarrow$ Renormalizable QFT: continuum limit

Universality $\leftrightarrow$ small number of relevant directions

Universality of the FP ensures that any discretization leads to the same continuum limit if the same properties concerning

- degrees of freedom
- locality
- symmetries

In asymptotically-free theories (QCD) the existence of a FP is warrantied by asymptotic freedom: can be proven in lattice perturbation theory!

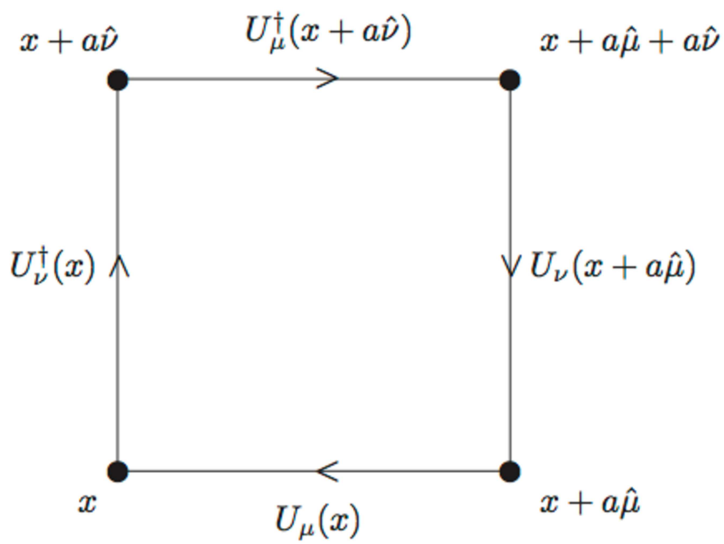
In general looking for scaling regions and FPs could be like looking for needles in a haystack because the dimension of bare coupling space is infinite

The hope is that a FP beyond perturbation theory also has a small number of relevant directions

Hierarchy problem: presence of relevant couplings: RGs pushes away from the FP $\rightarrow$ fine-tunning is needed

Triviality problem: the continuum limit is a free theory

Massive gauge fields on the lattice



$$U_\mu(x) \in SU(N)$$

$$U_\mu^\Omega(x) = \Omega^\dagger(x) U_\mu(x) \Omega(x+a\hat{\mu})$$

Any theory of gauge fields with non-gauge invariant interactions is equivalent to a gauge theory + charged scalar degrees of freedom:

$$\begin{aligned} Z &= \int dU \exp\left[-S_g[U] - S_{ni}[U]\right] = \int d\Omega \int dU \exp\left[-S_g[U] - S_{ni}[U]\right] \\ &= \int d\Omega \int dU^\Omega \exp\left[-S_g[U^\Omega] - S_{ni}[U^\Omega]\right] = \int d\Omega \int dU \exp\left[-S[U,\Omega]\right]. \end{aligned}$$

Exact gauge invariance:

$$S[U,\Omega] = S[U^{\Lambda^\dagger}, \Lambda\Omega]$$

In particular

Gauge theory with a mass term = Gauged non-linear σ model

$$S_m[U] = \beta \sum_{x,\mu,\nu} \mathrm{Retr}\,[1 - P(x,\mu,\nu)] - \frac{\kappa}{2} \sum_{x,\mu} \mathrm{tr}\left[U_\mu + U_\mu^\dagger\right],$$

$$\beta \equiv \frac{2N}{g_0^2}, \qquad \kappa \equiv \frac{\beta}{N}(ma)^2$$

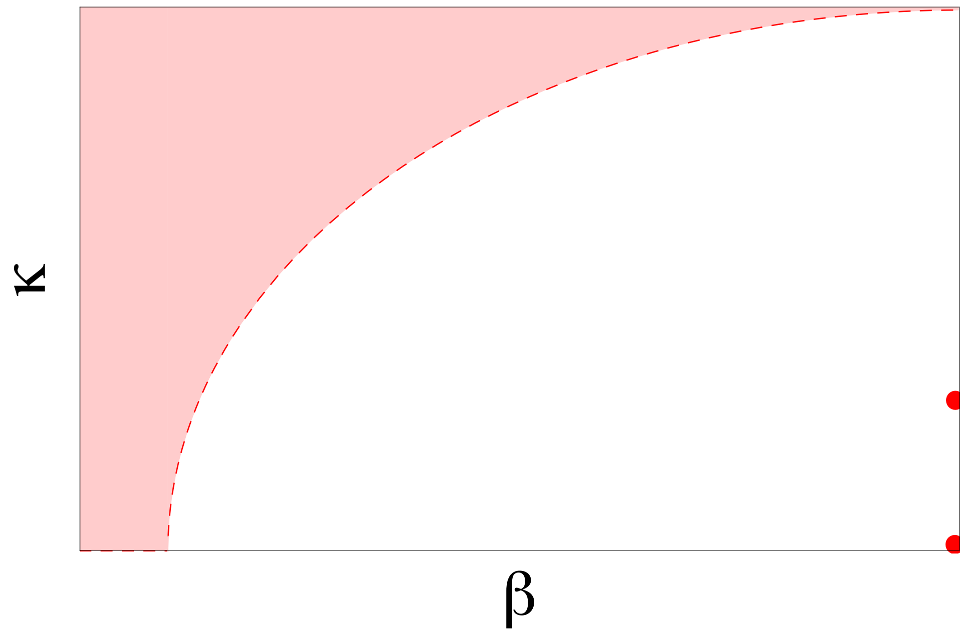
$$S[U,\Omega] \equiv \beta \sum_{x,\mu,\nu} \mathrm{Retr}\,[1 - P(x,\mu,\nu)] - \frac{\kappa}{2} \sum_{x,\mu} \mathrm{tr}\left[\Omega^\dagger(x) U_\mu(x) \Omega(x + a\hat{\mu}) + h.c.\right].$$

$$lim_{a\to 0}\; S[U,\Omega] = \frac{1}{4} W^a_{\mu\nu} W^a_{\mu\nu} + \frac{\kappa}{2} \mathrm{tr}[D_\mu \Omega^\dagger D_\mu \Omega]$$

For $N = 2$: this is also $SU(2)$+$\lambda\Phi^4$ in the limit $\lambda \to \infty$ (no fundamental Higgs)

Phase Diagram: on the sides

Lang, Rebbi, Virasoro; Osterwalder, Seiler; Fradkin, Shenker; Forster, Nielsen, Ninomiya



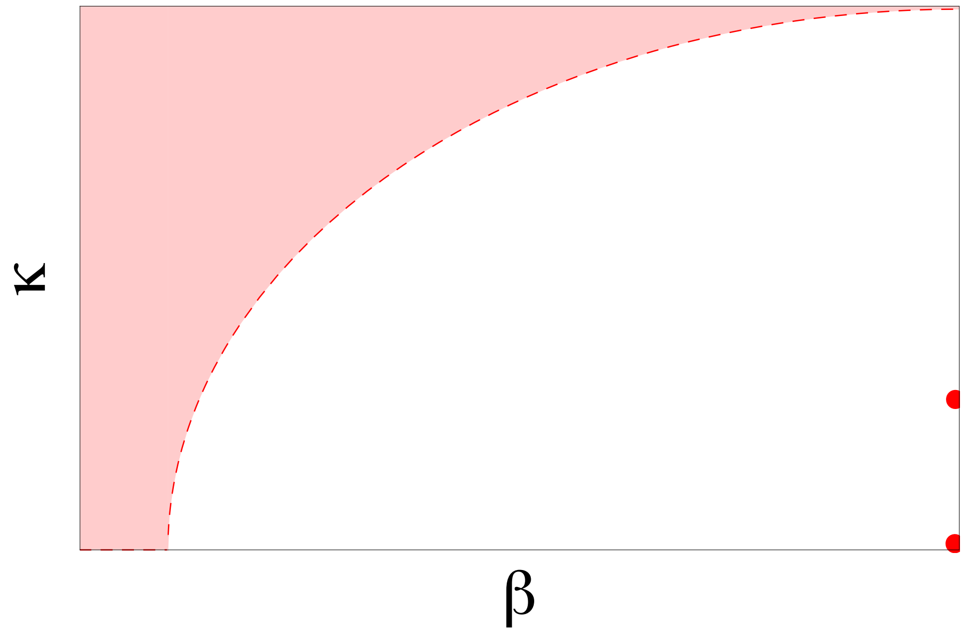
- $\kappa = 0$ axis: $SU(2)$ gauge theory, FP at $\beta_c = \infty$

- $\beta = \infty$ axis:non-linear σ-model, FP at κ_c: separates a disordered phase at $\kappa < \kappa_c$ from an ordered one at $\kappa > \kappa_c$ (Higgs or Goldstone phase)

  Perturbatively non-renormalizable, but lattice studies have shown that it is in the same universality class as the linear model (just one relevant parameter κ)

Phase Diagram: inside

Lang, Rebbi, Virasoro; Osterwalder, Seiler; Fradkin, Shenker; Forster, Nielsen, Ninomiya



- $\kappa < \kappa_{\rm min}$: pure gauge theory universality class (continuum limit at $\beta \to \beta_c$)

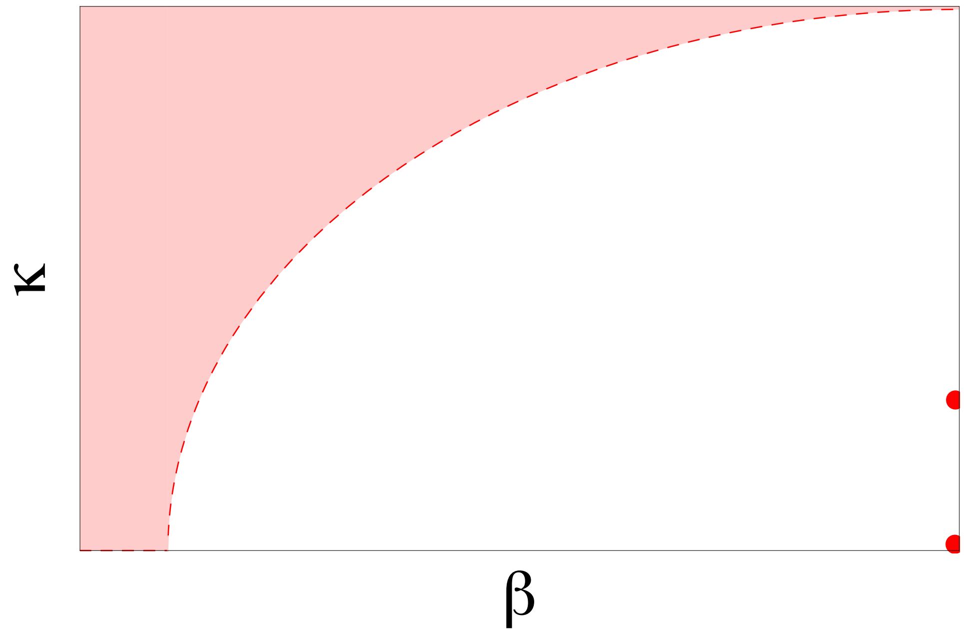
Foester, Nielsen, Ninomiya

$$\xi_\Omega / a = \text{finite}$$

Ω fields decouple as heavy dofs: the effect of S_{ni} can be reabsorbed in a renormalisation of the gauge-invariant operators.

Phase Diagram: inside

Lang, Rebbi, Virasoro; Osterwalder, Seiler; Fradkin, Shenker; Forster, Nielsen, Ninomiya



- confinement/Higgs phases: analytically connected in red region

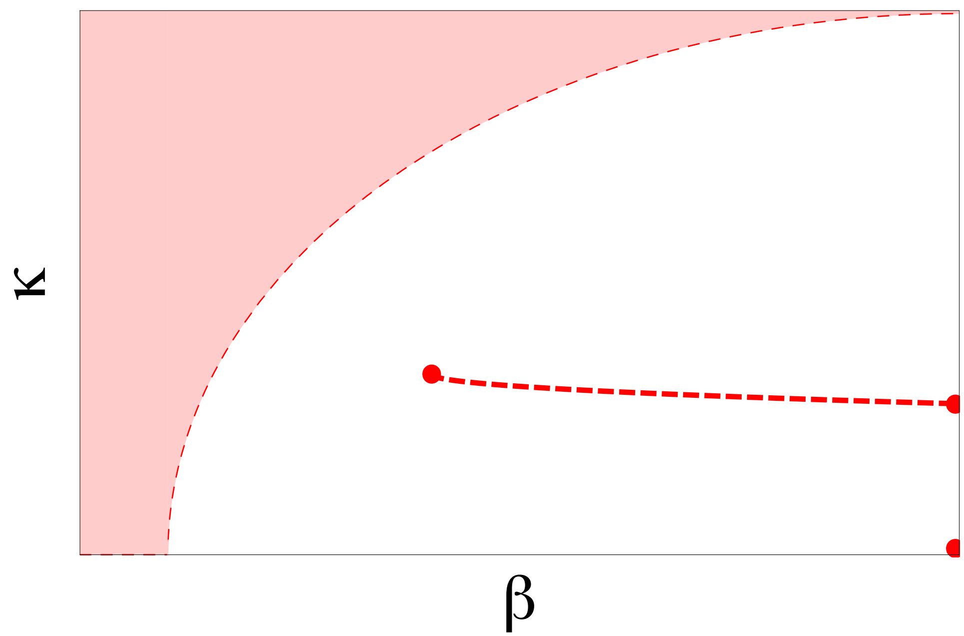
  The confinement/Higgs phase do not differ qualitatively at finite a, but they should differ in a scaling region

$$\xi_\Omega / a \to \infty \leftrightarrow \text{Higgs phase}$$

  Static charges can be screened by the Ω fields: static potential flattens at $r > r_0$ with $r_0/a \to \infty$.

Phase Diagram: inside

Langguth, Montvay, Weisz; Campos; Caudy, Greensite



Recent studies: line of first order phase transition end-point at $\beta_c \simeq 2.7$

Bonati, Cossu, D'Elia, Di Giacomo 2010

Is there a scaling region within the Higgs phase ? $\xi_\Omega / a \gg 1$?

If so what Wilsonian effective theory describes it ? What are the light dofs ? Is it non-perturbatively renormalizable (FP) ?

Strategy in this work: search for lines of constant physics and test scaling

Lattice perturbation Theory

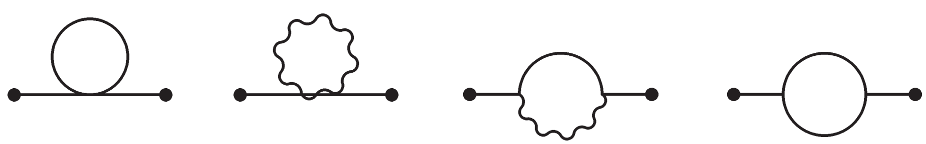
A perturbative expansion is possible as a low-momentum expansion (small g_0, large κ): $p^2 \ll \frac{\kappa}{a^2} = \frac{2m_W^2}{g_0^2} \equiv \frac{v^2}{2}$

(Equivalent to Chiral Perturbation Theory: natural cutoff $\sim 4\pi v = 4\pi \frac{\sqrt{2\kappa}}{a}$)

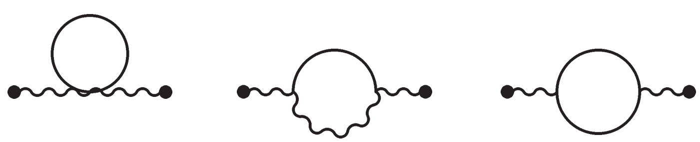
Appelquist, Bernard

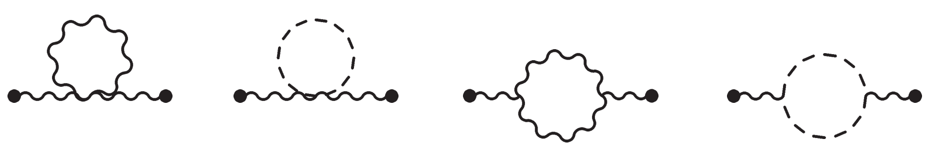
In the regime $m_W^2 \ll p^2 \ll \frac{2m_W^2}{g_0^2}$ with background field method (B_μ^a, ω^a) :

Lüscher, Weisz



$= -\frac{\kappa+\Delta\kappa}{4}\hat{p}^2\omega^a(-p)\omega^a(p)$





$=$

$$\left\{-\frac{1}{2(g^2+\Delta g^2)}\left[g_{\mu\nu}\hat{p}^2 - \hat{p}_\mu\hat{p}_\nu\right] - \frac{1}{4}(\kappa+\Delta\kappa)g_{\mu\nu}\right\} B_\mu^a(-p)B_\nu^a(p)$$

Up to corrections of $\mathcal{O}(m_W^2/p^2, p^2/\kappa)$:

$$\left(\frac{\Delta g^2}{g^4}\right) = \frac{N}{(4\pi)^2}\left(-\frac{29}{8}\ln p^2 + \frac{63}{9}\right) + N\left(\frac{7}{48}P_1 + \frac{29}{8}P_2 + \frac{1}{16}\right) - \frac{1}{8N}$$

$$\begin{aligned}\frac{\Delta\kappa}{\kappa} &= \frac{1}{\kappa}\left[\frac{1}{8N} - \frac{N}{16} - \frac{N}{2}P_1\right] \\ &- g^2\left[\left(\frac{5N}{32} - \frac{3}{16N}\right)P_1 + \frac{3N}{4}P_2 - \frac{N}{4}\frac{1}{(4\pi)^2} - \frac{3N}{4(4\pi)^2}\log(m_W^2)\right. \\ &+ \left.\frac{2N}{(4\pi)^2}F\left(\frac{m_W^2}{p^2}\right)\right],\end{aligned}$$

$$P_1 \equiv \int_{-\pi}^{\pi}\frac{d^4p}{(2\pi)^4}\frac{1}{\hat{p}^2} = 0.15493339...$$

$$P_2 \equiv \lim_{\mu\to 0}\left\{\frac{1}{(4\pi)^2}\log(\mu^2) + \int_{-\pi}^{\pi}\frac{d^4p}{(2\pi)^4}\frac{1}{(\hat{p}^2+\mu^2)^2}\right\} = 0.02401318....$$

$$F(x) \equiv 1 - \sqrt{1+4x}\ \mathrm{arccoth}\sqrt{1+4x}$$

Ungauged result $g = 0$: Shushpanov, Smilga

There is asymptotic freedom

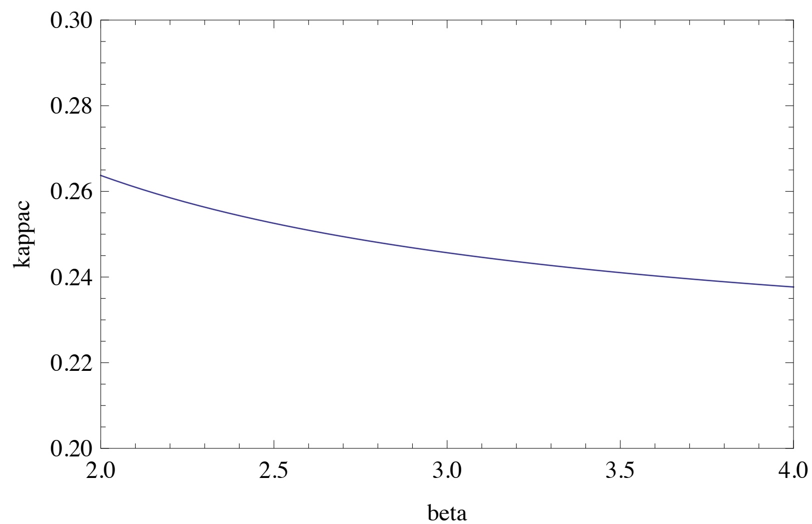
$$\beta(g) = -\frac{29N}{8(4\pi)^2} g^3 + \ldots,$$

Fabbrichesi et al

but with a different coefficient as in the pure gauge theory.

A continuum limit ? At this order requires:

$$\kappa + \Delta\kappa = 0$$



and a tunning order by order...

Wilsonian Effective Theory

Let us assume that a scaling region exists where $m_{\rm phys} a \to 0$ with the following properties:

- Asymptotic states are gauge singlets: confinement
- The lightest state has the quantum numbers of the W^a_μ boson (isovector) and is weakly coupled

Eg: interpolating field

$$V^a_\mu \equiv i\frac{\kappa}{2}{\rm Tr}[\Omega(x)^\dagger U_\mu(x)\Omega(x+\hat{\mu})T^a] + h.c.$$

Then the Wilsonian effective theory is itself a massive gauge theory up to effects of higher excited states

Exact global symmetry: Custodial Symmetry

The lattice action preserves an exact $SU(2)$ global symmetry:

$$\Omega(x) \quad \rightarrow \quad \Omega(x)\Lambda, \;\; \Lambda \in SU(N)$$

The corresponding conserved Noether currents are:

$$V_\mu^a \equiv i\frac{\kappa}{2}\mathrm{Tr}[\Omega(x)^\dagger U_\mu(x)\Omega(x+\hat{\mu})T^a] + h.c.$$

$$\hat{\partial}_\mu V_\mu^a = 0,$$

where $\hat{\partial}_\mu \Omega(x) \equiv \Omega(x+\hat{\mu}) - \Omega(x)$

The effective theory must preserve the global symmetry:

$$\begin{aligned} \mathcal{L}_{eff}(W) \quad &= \quad -\frac{1}{4}Z\; \partial_{[\mu,}W_{\nu]} \cdot \partial_{[\mu,}W_{\nu]} - \alpha\; W_\mu \times W_\nu \cdot \partial_\mu W_\nu \\ &- \quad Zm_W^2\; W_\mu \cdot W_\mu + \lambda(W_\mu \cdot W_\mu)^2 + \mu(W_\mu \cdot W_\nu)^2, \end{aligned}$$

Imposing that W transforms as a current under an infinitesimal symmetry transformation $\Lambda = e^{iT^a \epsilon^a(x)}$:

$$W_\mu \to \Lambda^\dagger W_\mu \Lambda + i\Lambda^\dagger \partial_\mu \Lambda$$

and that it is the conserved current of the global symmetry in the effective theory:

$$\frac{\partial \mathcal{L}_{eff}}{\partial \partial_\mu \epsilon^a(x)} \propto W_\mu^a(x),$$

then

$$\alpha = -4\lambda = 4\mu = Z.$$

while m_W^2 is not constrained. After canonically normalization:

$$\begin{aligned} \mathcal{L}_W &= -\frac{1}{4}\partial_{[\mu,}W_{\nu]} \cdot \partial_{[\mu,}W_{\nu]} - g\ W_\mu \times W_\nu \cdot \partial_\mu W_\nu - m_W^2\ W_\mu \cdot W_\mu \\ &- \frac{g^2}{4}\left[(W_\mu \cdot W_\mu)^2 - (W_\mu \cdot W_\nu)^2\right], \end{aligned}$$

with $g \equiv Z^{-1/2}$. This is a massive Yang-Mills theory!

Exact global symmetry: Ward Identities

For any operator O and a local infinitesimal rotation, $\Lambda(x) = e^{iT^a \epsilon^a(x)}$:

$$\langle -\delta S \; O \rangle + \langle \delta O \rangle = 0,$$

with

$$\delta_\epsilon S[U, \Omega] = -\sum_{x,\mu,a} V^a_\mu(x) \hat{\partial}_\mu \epsilon^a(x),$$

Case I: $O(y, z) \equiv V^a_\mu(y) V^b_\nu(z)$

$$\delta_\epsilon V^a_\mu(x) = -\epsilon^{abc} \left[\epsilon^b(x) \; V^c_\mu(x) - \frac{1}{2} \hat{\partial}_\mu \epsilon^b(x) \; V^c_\mu(x) \right] + \frac{1}{4} V^0_\mu(x) \hat{\partial}_\mu \epsilon^a(x),$$

where V^0_μ is a singlet under the global symmetry:

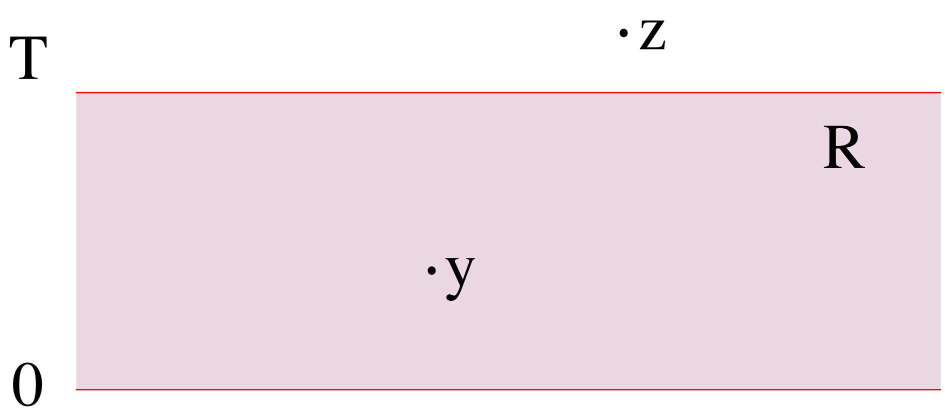
$$V^0_\mu(x) \equiv \frac{\kappa}{2} \mathrm{Tr}[\Omega(x)^\dagger U_\mu(x) \Omega(x + \hat{\mu}) + h.c.]$$

We consider

$$\epsilon^a(x) = \begin{cases} \epsilon^a, & x \in R \\ 0, & x \notin R \end{cases}$$

$y \in R$ (ie. $0 < y_0 < T$) while $z \notin R$, for example $z_0 > T$

Lüscher



The lattice WI implies

$$\epsilon_{abc} \sum_{\mathbf{x}} \langle \left(V_0^c(T, \mathbf{x}) - V_0^c(0, \mathbf{x}) \right) V_\mu^a(y) V_\nu^b(z) \rangle = 2 \langle V_\mu^d(y) V_\nu^d(z) \rangle,$$

Matching to the effective theory

$$V_\mu^a = Z_W^{1/2} W_\mu^a \equiv m_W F_W W_\mu^a$$

and evaluating the three-point function to LO in perturbation theory:

$$\frac{1}{\sqrt{Z_W}} = \frac{g}{m_W^2} \to g = \frac{m_W}{F_W}$$

The effective gauge coupling g can be extracted from the two-point function:

$$\lim_{T\to\infty} \sum_{\mathbf{x}} \langle V_i^a(T,\mathbf{x}) V_i^a(0,\mathbf{0})\rangle = \frac{m_W F_W^2}{2} e^{-m_W T}$$

Case II: $O \to V_\mu^b(y) V_\nu^c(z) V_\sigma^d(u)$

with $a = b = c = d$ and $y, z, u \notin R$

$$\lambda = -\mu.$$

Case III: $O \to V_\mu^b(y) V_\nu^c(z) V_\sigma^d(u)$ with $a = b \neq c = d$

$$\lambda Z_W^{-2} = \frac{-g^2}{4} \to .$$

Global invariance $\Rightarrow$ local invariance in the effective theory of conserved currents

Wilsonian effective theory of massive gauge bosons + Higgs ?

Possibly other light particles will exist in the spectrum, in particular a singlet scalar such as

$$H(x) \leftrightarrow V_\mu^0(x) \equiv \frac{\kappa}{2}\mathrm{Tr}[\Omega(x)^\dagger U_\mu(x)\Omega(x+\hat{\mu}) + h.c.]$$

If such a light particle remains in the spectrum how is the effective theory modified ?

Global symmetry allows the following couplings:

$$\mathcal{L}_H = \frac{1}{2}\partial_\mu H \partial_\mu H - V(H) - \lambda_{HWW} H W_\mu \cdot W_\mu - \lambda_{HHWW} H^2 W_\mu \cdot W_\mu,$$

The WI in this case implies the following matching:

$$V_\mu^a \to W_\mu^a + 2\frac{\lambda_{HWW}}{m_W^2} H W_\mu^a + ...$$

Symmetry does not seem to require the couplings to be those in the SM , but we know unitarity does...

Scaling Studies

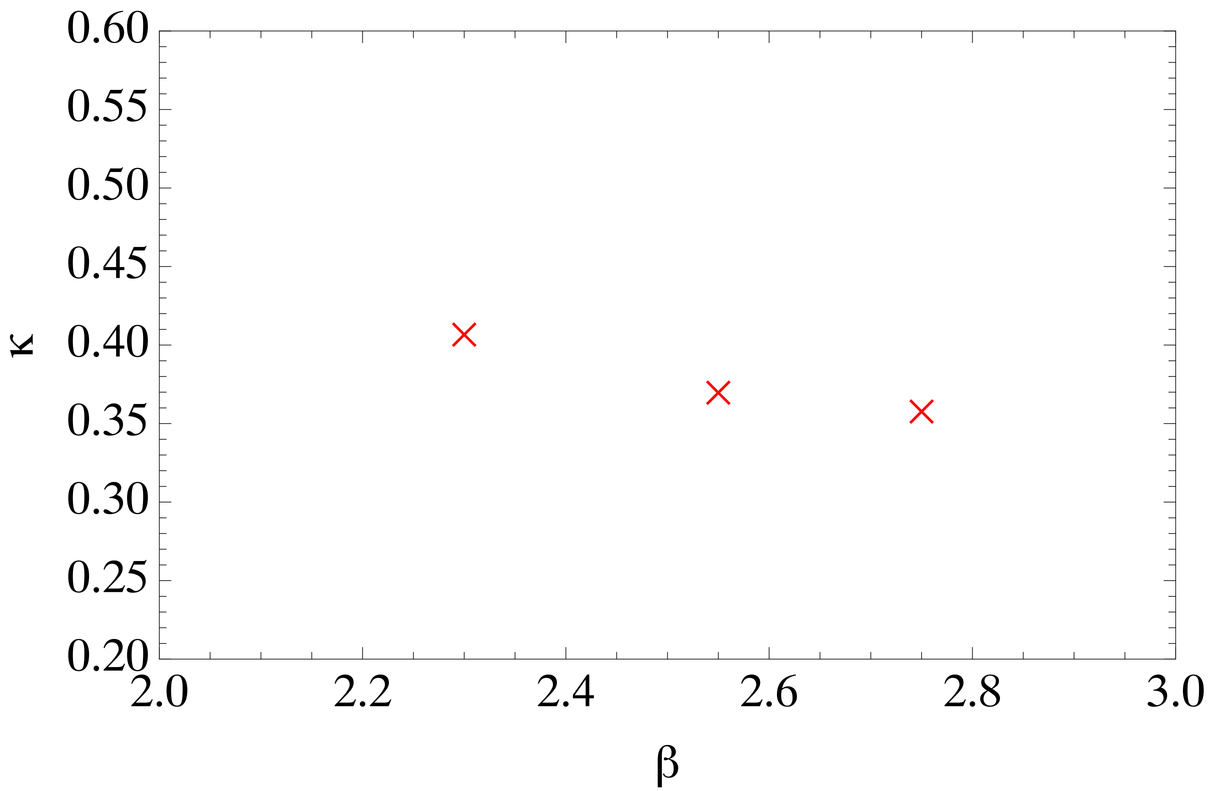
Basic strategy

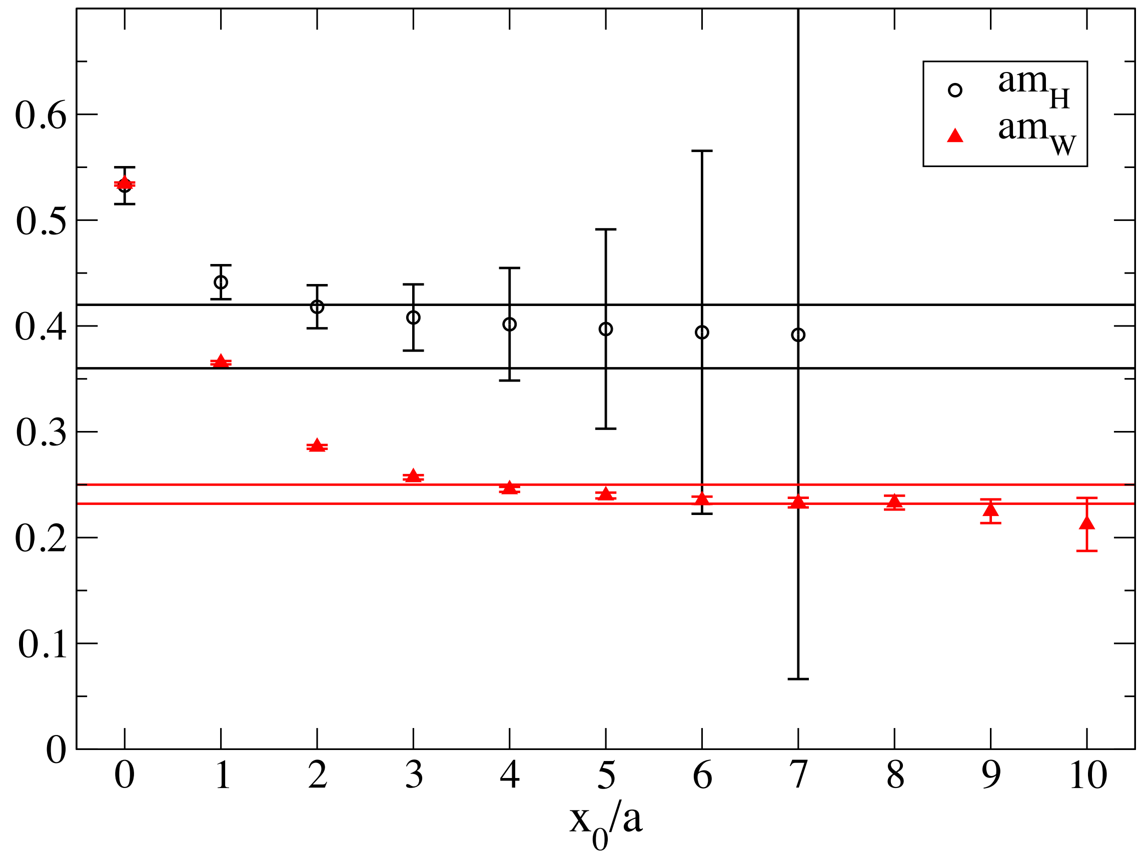
- We start with the benchmark point $\beta = 2.3$ and $L = 16$. We tune κ to reach $m_H/m_W \simeq 1.4$
- We increase β and L to keep $m_W L \geq 5$ and perform the same tunning in κ to keep m_H/m_W fixed.
- We test scaling of the observables: m_W, m_H, F_W and static potential.

$$
\begin{aligned}
V(t) &= -\kappa^2 \sum_{\vec{x}} \sum_{k=1}^{3} \sum_{a} \langle \mathrm{Tr}[U_k(x)T^a]\mathrm{Tr}[U_k(0)T^a]\rangle \ , \\
S(t) &= \kappa^2 \sum_{\vec{x}} \langle \sum_{\mu} \mathrm{Tr}[U_\mu(x)] \sum_{\nu} \mathrm{Tr}[U_\nu(0)]\rangle \ , \\
V_{stat}(r) &= -\frac{1}{T} \log C_{PP}(r)_{\mathrm{connected}}.
\end{aligned}
$$

β	κ_{ref}	$L^3 \times T$	am_H	am_W	aF_W	$N_{\mathrm{meas}}/10^6$
2.3	0.405	$16^3 \times 16$	0.65(2)	0.455(5)	0.146(2)	5.4
2.55	0.368	$24^3 \times 24$	0.39(3)	0.241(9)	0.081(2)	2.8
2.75	0.356	$36^3 \times 36$	0.30(4)	0.174(6)	0.060(2)	2.0

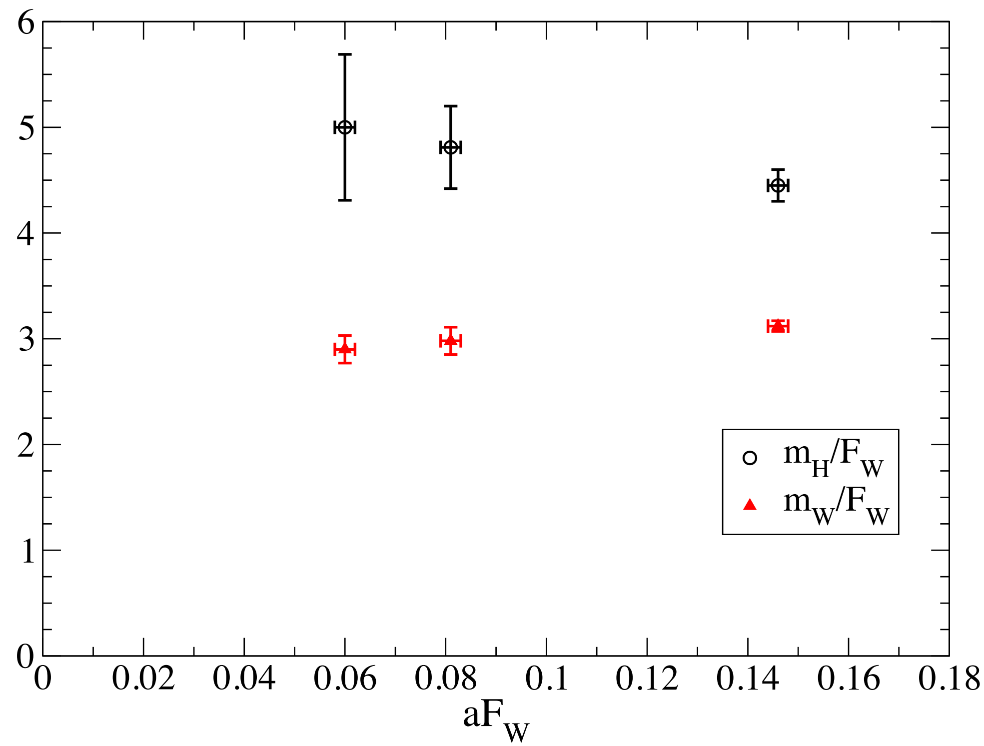
Table 1: Lattice parameters and estimates for am_H, am_W and aF_W.





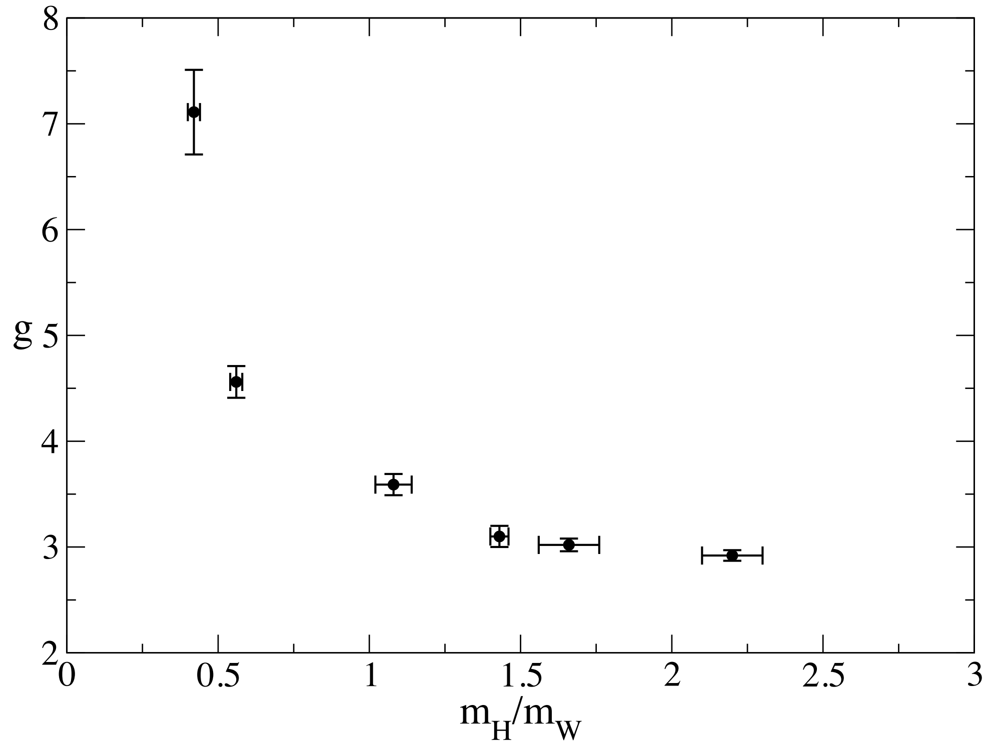
am$_H$
am$_W$
0
0.1
0.2
0.3
0.4
0.5
0.6
0
1
2
3
4
5
6
7
8
9
10
x_0/a

	am_{H}	am_{W}	aF_{W}
$\frac{\beta=2.3}{\beta=2.55}$	1.67(14)	1.89(7)	1.80(5)
$\frac{\beta=2.55}{\beta=2.75}$	1.30(20)	1.38(7)	1.35(6)

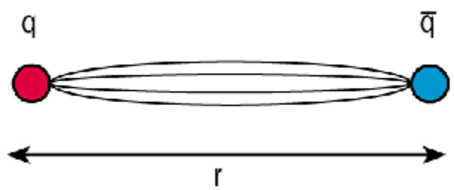


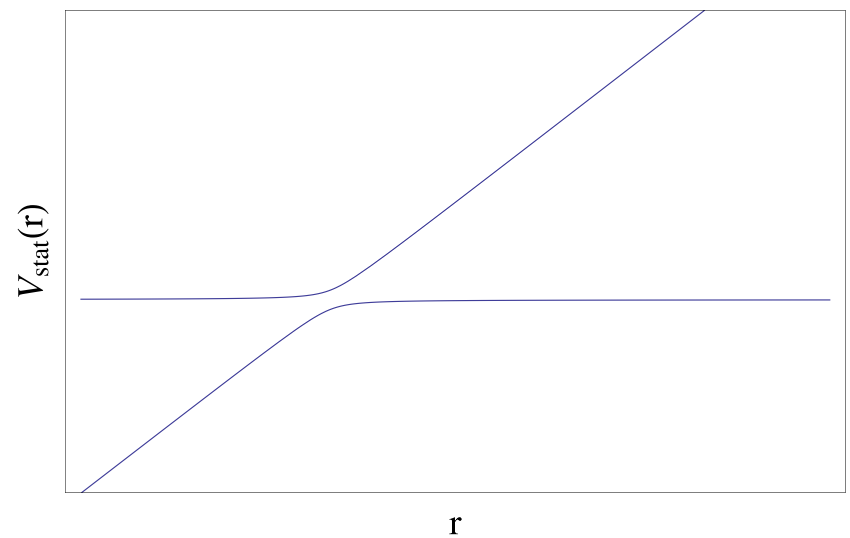
Very large g...

Smaller couplings for different m_H/m_W ?



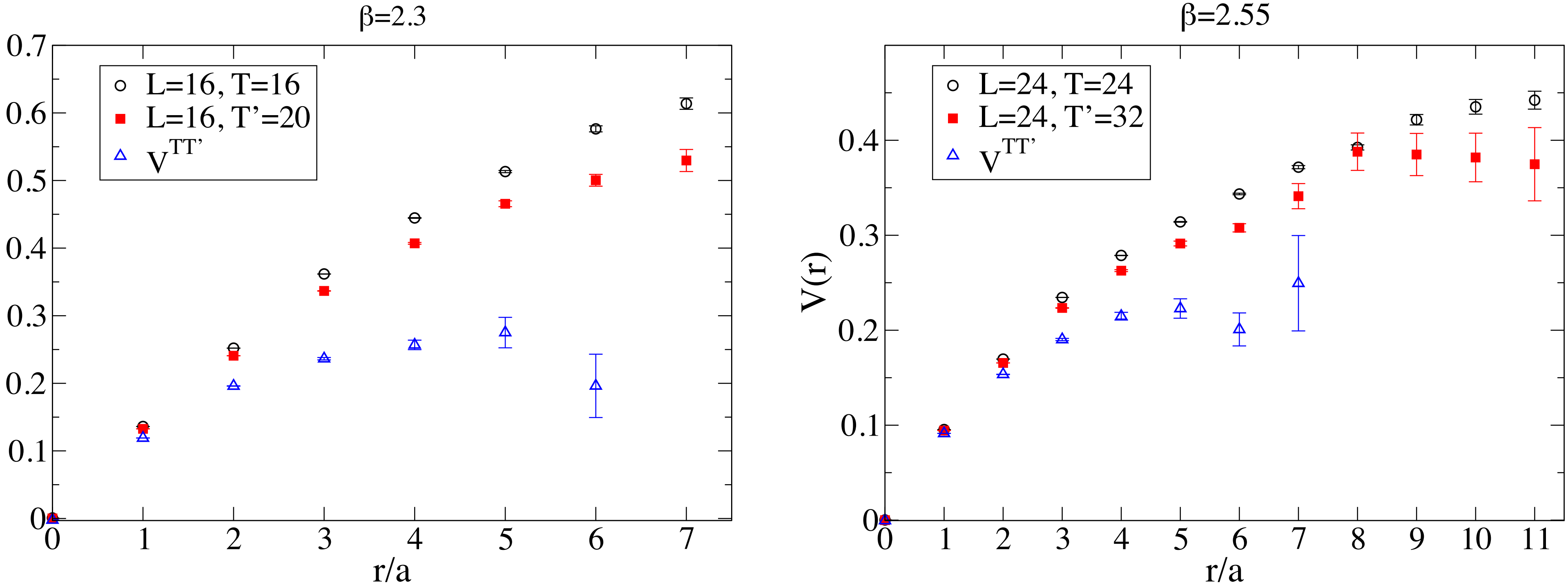
Flattening of $V_{stat}(r)$ for large r expected from string breaking due to Ω-states: mixing of stringy/static-Ω:





The overlap of Polyakov loop on the ground state depends on r

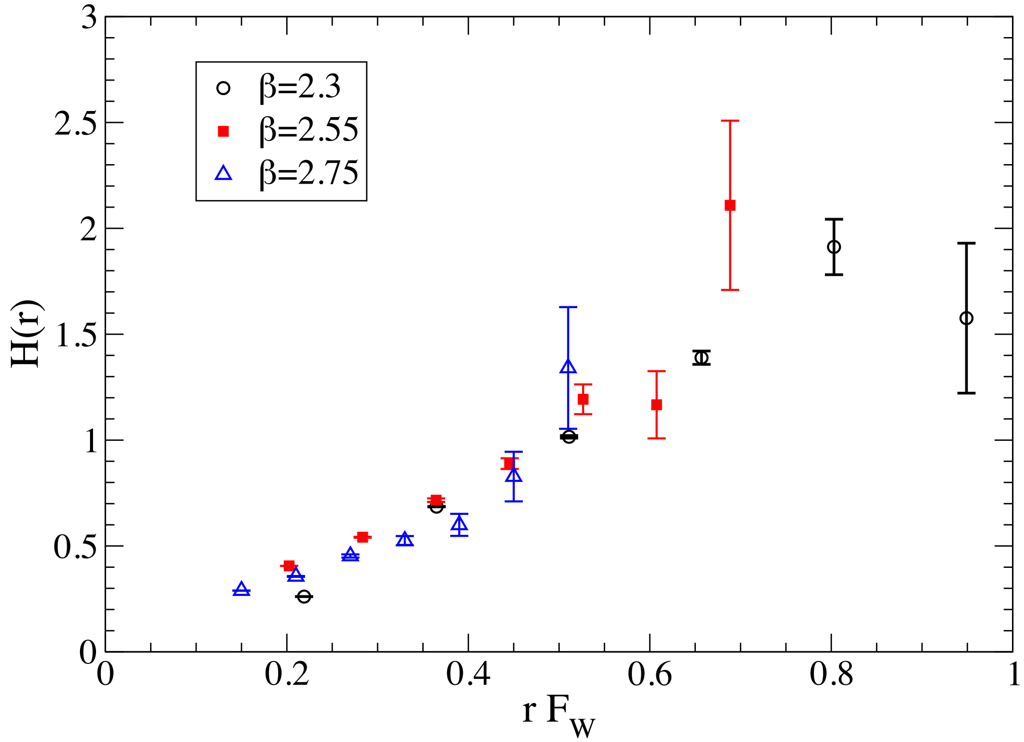
$$C_{PP}(r,T) = w_0(r) e^{-V_{stat}(r)T}, \qquad V^{TT'} = -\frac{1}{T-T'} \log\left(\frac{C_{PP}(r,T)}{C_{PP}(r,T')}\right)$$



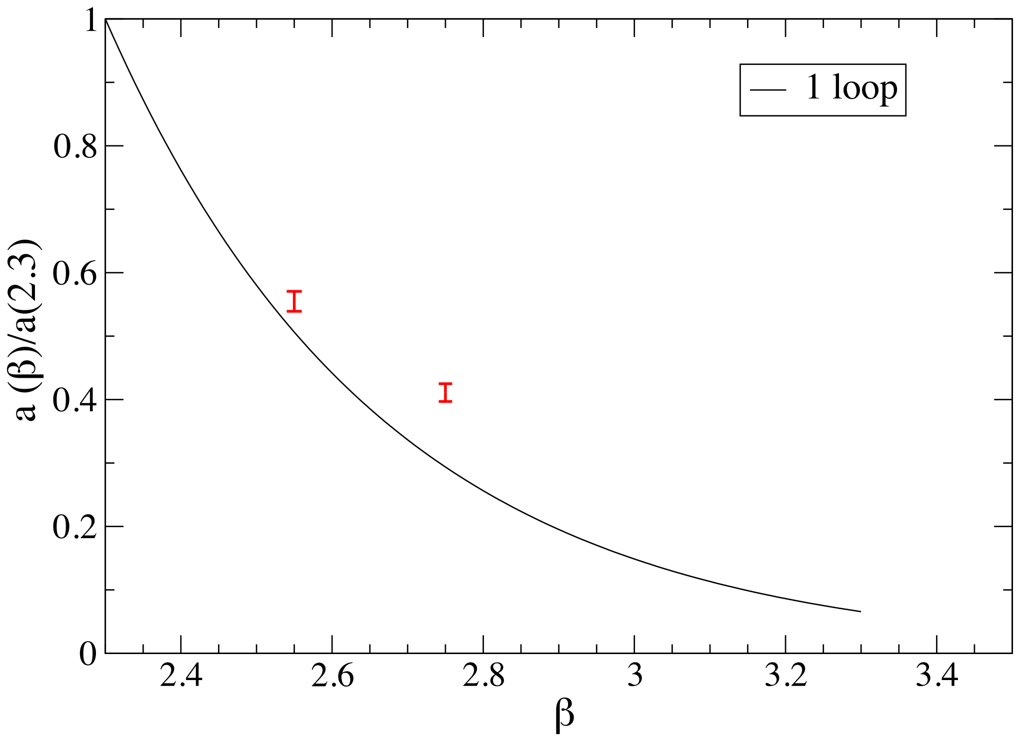
β=2.3
L=16, T=16
L=16, T'=20
V^TT'
r/a
β=2.55
L=24, T=24
L=24, T'=32
V^TT'
V(r)
r/a

Scaling of the static potential

$$H(r) \equiv r^2 \frac{\partial V_{stat}(r)}{\partial r}$$



Comparison with one-loop scaling



Perturbation theory is not the whole story

Conclusions and Outlook

- A gauged non-linear sigma model might be the simplest model for dynamical EWSB: the scalar dofs are part of the gauge field
- This is a non-trivial strongly coupled theory that needs to be understood non-perturbatively: it might be renormalizable and more universal than we think
- Global symmetries indicate that if a continuum limit/scaling region exists it could look very similar to a SSB gauge theory: massive gauge fields and light scalar
- Old studies can be very significantly improved with new methods and algorithms developed for QCD
- Several improvements under way: improve signal for scalar, measure vector-scalar couplings, measure Wilson loops to solve mixing problem, excited states, universality test