The Phenomenological Viability of
Top Condensation

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Abstract

We discuss how the full dynamics of top condensation models can modify the relations between the physical top mass, the amount of custodial $SU(2)$ violation and the weak gauge boson masses. It is emphasized that it is possible to get phenomenologically acceptable relations between $\Delta \rho$, $m_t$ and $M_W$ and that in addition the scale of new physics can be chosen to be $\mathcal{O}(TeV)$ such that a fine–tuning problem is avoided.
The formation of a $\bar{t}t$ condensate by some “pairing force” could be responsible for a dynamical breaking of the electro–weak symmetry [1, 2]. A theory creating such a condensate would naturally explain a heavy top mass, would be very helpful to avoid Flavour Changing Neutral Current (FCNC) problems and would be very attractive due to its economy. Simple initial realizations of top condensation were based on effective Nambu–Jona-Lasinio (NJL) models with non–renormalizable four–fermion interactions. This led however to discussions about higher dimensional operators [3] which depend crucially on how “effective” or “fundamental” the four–fermion interactions are. Subsequently fundamental four–fermion theories where proposed [4] while other authors justify an effective NJL description of strongly coupled broken gauge theories [5, 6, 7, 8, 9].

Independently of such questions we study here the phenomenological viability of top condensation ideas by assuming essentially only that we know the solution $\Sigma_t(p^2)$ of some relevant Schwinger–Dyson (“gap”) equation for the dynamically generated top mass. Thus we pretend to know the electro–weak symmetry breaking top propagator to be

$$S_t(p^2) = \frac{i}{p^2 - \Sigma_t(p^2)},$$

with the (pole) top mass $m_t = \Sigma_t(m_t^2)$. All other quarks and leptons are assumed to be massless. Without specifying the gap equation we assume furthermore that for the theory under consideration $\Sigma_t(p^2) \rightarrow 0$ and that there is only one unique solution for $m_t$.

The breaking of the electro–weak symmetry (i.e. $\Sigma_t \neq 0$) is assumed to be the result of unspecified new strong forces acting only on the known quarks and leptons and especially on the $t–b$ doublet. The emergence of a top condensate breaks global symmetries and the resulting Goldstone Bosons are “eaten” in a dynamical Higgs mechanism such that $W$ and $Z$ become massive. Presumably such a theory does not change significantly if the weak $U(1)_Y$ coupling $g_1$ is sent to zero. In the limit $g_1 = 0$ the corrections which give mass to the $W^\pm$ and $W_\pm$ propagators must be induced by the fermions which are representations under both $SU(2)_L$ and the new strong force. We should therefore study the contributions of $\Sigma_t$ to the vacuum polarizations of the $W$ and $Z$ propagators. In an expansion in powers of $g_2^2$ the leading contribution is given by diagrams which connect the $W_\pm$ or $W_3$ line to a fermion pair from both sides. There are two ways how these four internal fermion lines can be connected: By inserting twice the full fermionic propagators or by inserting once the full four–fermion Kernel of the new, strong interaction. Note that in leading order $g_2^2$ the fermion propagators and the Kernel do not contain any electro–weak gauge boson propagation themself since this would cost at least an extra power of $g_2^2$. Insertions of fermionic vacuum polarizations into higher order electro–weak loop diagrams,

\[1\] This is e.g. justified for asymptotically free theories where chiral symmetry breaking disappears as $p^2 \rightarrow \infty$.

\[2\] Indirectly (via vacuum alignment) a small $U(1)_Y$ coupling could be very important such that $g_1 = 0$ should be understood as the result of the limiting procedure $g_1 \rightarrow 0$. 


for example, are suppressed by corresponding powers of $g_2^2$. Thus in leading order $g_2^2$, but exact in the new strong coupling, the $W$ propagator is graphically represented by Fig. 1. The first contribution is a generalization of the leading Standard Model diagrams with hard fermion masses replaced by $\Sigma$’s, i.e. the sum of all one particle irreducible diagrams which contribute to the dynamically generated fermion masses. The second contribution contains the exact Kernel $K$ of the strong forces responsible for condensation and it is useless to expand this Kernel perturbatively in powers of the coupling constant of the new strong force. The Goldstone theorem tells us however that the Kernel must contain poles of massless Goldstone Bosons due to the breaking of global symmetries by the fermionic condensates. This is symbolically expressed by the second line of Fig. 1, where $\tilde{K}$ does not contain any further poles of massless particles. But $\tilde{K}$ may (and typically will) contain all sort of massive bound states which could e.g. be vectors, Higgs–like scalars etc. in all possible channels.

The Goldstone Boson contributions shown in the second line of Fig. 1 were used by Pagels and Stokar [11] to obtain a relation between the $\Sigma$’s and the Goldstone Boson decay constants. Their derivation uses the fact that only the Goldstone Bosons contribute a term proportional $p_\mu p_\nu/p^2$ to the $W$ polarization at vanishing external momentum, but this method ignores possible contributions from $\tilde{K}$ which enter indirectly via the use of Ward identities. The $p_\mu p_\nu/p^2$ contributions to $\Pi_{\mu\nu}$ are balanced (up to small corrections from $\tilde{K}$) by $g_{\mu\nu}$ terms created by the first diagram on the rhs of Fig. 1. Following ref. [10] we derive a relation between the $\Sigma$’s and the Goldstone Boson decay constants from these $g_{\mu\nu}$ terms. The result can be compared with the Pagels–Stokar relation and we will see that contributions from $\tilde{K}$ are significantly suppressed. Let us therefore work with rescaled fields such that gauge couplings appear in the kinetic terms of the gauge boson Lagrangian like $(-1/4 g^2) (W_{\mu\nu})^2$. Since we do not include any propagating $W$ bosons we need not gauge fix at this stage. The inverse $W$ propagator can be written as

$$\frac{1}{g_2^2} D_{W,\mu\nu}(p^2) = \frac{1}{g_2^2} \left( -g_{\mu\nu} + \frac{p_\mu p_\nu}{p^2} \right) p^2 - \Pi_{\mu\nu}(p^2),$$

(2)

with the polarization tensor $\Pi_{\mu\nu}(p^2) = (-g_{\mu\nu} p^2 + p_\mu p_\nu) \Pi(p^2)$. At vanishing external momentum the first fermion loop on the rhs of Fig. 1 contributes to $\Pi_{\mu\nu}$

$$\Pi_{\mu\nu} = -i Z^2 N_c \int \frac{d^4 k}{(2\pi)^d} \frac{\text{Tr} [\Gamma_\mu(k + \Sigma_1(k))\Gamma_\nu(k + \Sigma_2(k))]}{(k^2 - \Sigma_1(k)^2)(k^2 - \Sigma_2(k)^2)},$$

(3)

where $N_c$ is the number of colors, $Z^{-1} = \sqrt{2}, 2$ in the charged and neutral channel, respectively, $\Gamma_\alpha = \left( \frac{1-\gamma_5}{2} \right) \gamma_\alpha$, and $+i\epsilon$ is generally implied in the denominator. In the neutral channel we get corrections from $\Omega$ (i.e. $\Sigma_1 = \Sigma_2 = \Sigma_t$), $\Omega_b$ (i.e. $\Sigma_1 = \Sigma_2 = \Sigma_b \equiv 0$) and in the charged channel contributes only $\Omega_b$ or $\Omega_t$ (i.e. $\Sigma_1 = \Sigma_t, \Sigma_2 = \Sigma_b \equiv 0$). By

\[3\]Which are essential for a gauge invariant dynamical Higgs mechanism.
naive power counting eq. (3) has quadratic and logarithmic divergences. Since we assumed 
\( \Sigma_i(p^2) \to 0 \) for the top quark and all other fermions we find that the divergences of 
\( \Pi_{\mu\nu}(p^2) \) are identical to those calculated for \( \Sigma_i \equiv 0 \). It makes therefore sense to split
\( \Pi_{\mu\nu}(p^2) = \Pi_{\mu\nu}^0(p^2) + \Delta \Pi_{\mu\nu}(p^2) \) where \( \Pi_{\mu\nu}^0 \) is defined as \( \Pi_{\mu\nu} \) for \( \Sigma_i \equiv 0 \). \( \Pi_{\mu\nu}^0 \) is then an
uninteresting \( \Sigma_i \) independent constant which contains all divergences and needs renormal-
ization. Contrary the interesting \( \Sigma_i \) dependent piece \( \Delta \Pi_{\mu\nu} = \Pi_{\mu\nu} - \Pi_{\mu\nu}^0 \) is finite, even
when the external momentum is sent to zero. Thus

\[
\Delta \Pi_{\mu\nu} = -iZ^2N_c \int \frac{dk}{(2\pi)^4} \left\{ \frac{Tr \left[ \Gamma_{\mu}(k + \Sigma_1)\Gamma_\nu(k + \Sigma_2) \right]}{(k^2 - \Sigma_1^2)(k^2 - \Sigma_2^2)} - \frac{Tr \left[ \Gamma_{\mu}\Gamma_\nu \right]}{k^4} \right\}
\]

(4)

\[
= -g_{\mu\nu} \frac{Z^2N_c}{(4\pi)^2} \int_0^\infty dk^2 \frac{k^2(\Sigma_1^2 + \Sigma_2^2) - 4\Sigma_1^2\Sigma_2^2}{(k^2 - \Sigma_1^2)(k^2 - \Sigma_2^2)}
\]

(5)

where angular integration was performed in Euclidean space and subsequently continued
back to Minkowski space. Under the integral one has as usual \( Tr \left[ \Gamma_{\mu}\Gamma_\nu \right] = -g_{\mu\nu}k^2 \)
and \( Tr \left[ \Gamma_{\mu}\Gamma_\nu \right] = 0 \). Note that this separation procedure for \( \Delta \Pi_{\mu\nu} \) does not spoil
gauge invariance.

The Goldstone Boson decay constants \( F_i^2 \) are the poles of \( \Pi(p^2) \) at vanishing external
momentum. For our definition of \( \Pi_{\mu\nu} \) we find that \( F_i^2 \) is identical to eq. (3) without
the factor \(-g_{\mu\nu}\). Using \( Z \) for the charged and neutral channel one finds

\[
F_3^2 = \frac{N_c}{32\pi^2} \int_0^\infty dk^2 \frac{\Sigma_i^4}{(k^2 - \Sigma_i^2)^2}, \quad F_3^2 = \frac{N_c}{32\pi^2} \int_0^\infty dk^2 \frac{k^2\Sigma_i^2 - \Sigma_i^4}{(k^2 - \Sigma_i^2)^2}
\]

(6)

such that

\[
F_3^2 - F_3^2 = \frac{N_c}{64\pi^2} \int_0^\infty dk^2 \frac{\Sigma_i^4}{(k^2 - \Sigma_i^2)^2}.
\]

(7)

Eq. (3) for \( F_3^2 \) is equivalent to the result obtained by Pagels and Stokar [11] from the
\( q_\mu q_\nu/q^2 \) contributions of Goldstone Bosons to \( \Pi_{\mu\nu} \). The result for the neutral channel in
eq (3) looks however somewhat different. But by using the integral identity

\[
\int_0^\infty dx \frac{x^2f(x)' - f(x)^2}{(x - f(x))^2} = f(\infty),
\]

(8)

for \( x = k^2 \) and \( f = \Sigma_i^2 \) we can rewrite eq. (3) into

\[
F_3^2 = \frac{N_c}{32\pi^2} \int_0^\infty dk^2 \frac{k^2\Sigma_i^2 - k^2\Sigma_i^4}{(k^2 - \Sigma_i^2)^2},
\]

(9)

where \( \Sigma_i' = d\Sigma_i/dk^2 \). Even though this looks now formally similar to the Pagels-Stokar
result it differs by a factor 2 in front of the derivative term in the nominator of eq. (3). This
difference may appear less important, but we will see that in the limit of a hard top mass our method produces the correct $\rho$-parameter, while the Pagels–Stokar result produces 3/2 times the correct answer. Additionally our expression leads also to a better numerical estimate of $f_\pi$ if we follow the methods of ref. \[11\].

The $\rho$-parameter \[12\] is defined as $\rho := \frac{F^2_\pm}{F^2_\mp}$ which can now be written as

$$\rho = 1 + \Delta \rho = \frac{F^2_\pm}{F^2_\mp} = \left( 1 + \frac{(F^2_\mp - F^2_\pm)}{F^2_\pm} \right)^{-1} \simeq 1 - 2 \frac{(F^2_\pm - F^2_\mp)}{v^2},$$

and from eq. \(7\) we find the contribution of the $t - b$ doublet

$$\Delta \rho = \frac{-N_c}{32\pi^2v^2} \int_0^\infty dk^2 \frac{\Sigma^4_t}{(k^2 - \Sigma^2_t)^2},$$

where we used $F^2_\pm = v^2/2$ with $v \simeq 175$ GeV in the denominator. Model independent parametrizations of radiative corrections parametrize the information contained in $\Delta \rho$ essentially in the variables $T$ \[13\] or $\epsilon_1$ \[14\].

With the expressions for $\Delta \rho$ in eq. \(11\) and $F^2_i$ in eq. \(8\) we can calculate for given $\Sigma_t(p^2) \xrightarrow{p^2 \to \infty} 0$ three independent observable quantities which are one of the weak gauge boson masses (either $M^2_W = g^2_2 F^2_\pm$ or $M^2_Z = (g^2_1 + g^2_2) F^2_3$), $\Delta \rho$ and furthermore the physical top mass $m_t$. These three quantities are dominated by different momenta and therefore $\Sigma \neq constant$ leads to a different relation than a constant, i.e. hard mass. It is instructive to look at the degree of convergence of the involved integrals. The Goldstone Boson decay constants $F^2_i$ are formally log. divergent, but are finite with our assumption on $\Sigma_i(p^2)$. In that case renormalization is not needed, but due to the formal log. divergence $\Sigma_t$ contributes with equal weight at all momentum scales. In other words, the magnitude of $F^2_i$ depends crucially on the high energy tail of $\Sigma_t$. The difference $F^2_\pm - F^2_3$ has better convergence properties and is always finite, even for $\Sigma_i(p^2) = constant$. This implies that $\Delta \rho$ is finite, as it should be, and it is most sensitive to infrared scales somewhat above $m_t$. Finally $m_t$ is of course only sensitive to one point, namely $m_t = \Sigma_t(m_t^2)$.

We would like to study now corrections in the relation between $m_t$, $M_W$ and $\Delta \rho$ when $\Sigma_t$ is the solution of a hypothetical Schwinger–Dyson equation which deviates from $\Sigma_t = m_t = constant$. First we would like to see if the correct Standard Model result emerges for a $t - b$ doublet when $\Sigma_t \to m_t$. Therefore we set

$$\Sigma_t(p^2) = m_t \Theta(\Lambda^2 - p^2),$$

and ignore again the $b$ quark mass. From eq. \(11\) we obtain for our ansatz

$$\Delta \rho = \frac{N_c m_t^2}{32\pi^2v^2} \left( \frac{1}{1 - m_t^2/\Lambda^2} \right) \xrightarrow{\Lambda \to \infty} \Delta \rho^{SM} = \frac{N_c \alpha_{em}}{16\pi \sin^2 \theta_W \cos^2 \theta_W} \frac{m_t^2}{M^2_Z}. \quad (13)$$
Note that in the limit $\Lambda \rightarrow \infty$ (i.e. a hard, constant top mass) we obtain correctly the leading Standard Model value while the Pagels–Stokar relation would produce $3/2$ times the Standard Model result. For finite $\Lambda$ eq. (13) describes furthermore the modification of the Standard Model result due to a physical high energy momentum cutoff. Such a cutoff makes $\Delta \rho$ a little bit more positive than in the Standard Model which implies for a fixed experimental value of $\Delta \rho$ a lower top mass prediction. From eq. (6) it is in addition possible to determine $M_W$ for the ansatz eq. (12)

$$M_W^2 = g_2^2 F_\pm^2 = \frac{g_2^2 N_c}{32\pi^2} \int_0^\Lambda dk^2 \frac{\Sigma_t^2}{k^2 - \Sigma_t^2} = \frac{g_2^2 N_c}{32\pi^2} m_t^2 \ln \left( \frac{\Lambda^2 - m_t^2}{m_t^2} \right). \quad (14)$$

Taking as experimental input $M_W = 80.14 \pm 0.27$ GeV, $\Delta \rho = 0.005 \pm 0.008$, $\alpha_{em}^{-1}(M_Z^2) = 127.8 \pm 0.1$ and $\sin^2 \theta_{eff}^f(M_Z^2) = 0.2318 \pm 0.0007$ we plot in Fig. 2 the two central top mass values resulting from eqs. (13) and (14) as a function of $\Lambda$ (dashed lines).

The ansatz eq. (12) can be viewed as the result of a Nambu–Jona-Lasinio (NJL) gap equation of top condensation as for example in the model of Bardeen, Hill and Lindner (BHL) [1]. In fact a NJL gap equation is the simplest conceivable Schwinger–Dyson equation where $\Sigma_t$ is forced to be a constant. Fig. 2 shows clearly that ultra high values of $\Lambda$ and the experimental errors are required to get the two top mass values in agreement. For such high $\Lambda$ the effective Lagrangian is valid for many orders of magnitude which led in the BHL analysis to the so–called “renormalization group improvement”. This means in the current language that $\Sigma_t = constant$ is replaced by $\Sigma_t = g_t(p^2) v$, where $v = 175$ GeV and $g_t(p^2)$ is the solution of the one–loop renormalization group equation. In BHL the predicted top mass is then the “effective fixedpoint” of the renormalization group flow. The same result could be seen in eq. (14) since the effective fixedpoint dictates the shape of $\Sigma(p^2)$ for many orders of magnitude. The BHL scenario has however phenomenological problems. First the very high value of $\Lambda$ is nothing else then the old hierarchy problem which appears now as a fine–tuning of the four–fermion coupling $G$. Furthermore the infrared fixedpoint prediction is higher than the dashed curve resulting from eq. (14) which is shown in Fig. 2 and has (within newest experimental errors) no intersection with the line resulting from eq. (13). Thus this simplest scenario seems unacceptable even for very high values of $\Lambda$.

Remembering that $\Delta \rho$ and $M_W$ are sensitive to details of $\Sigma_t$ in a different way we should ask ourselves if the above problems can be solved by modifications of the solution $\Sigma_t(p^2)$. The answer is of course yes, and we illustrate now the two most important type of changes: The addition of a slowly falling tail and/or the addition of a “bump” somewhat above $m_t$.

First we consider a very rough ansatz for a “bump” between $\Lambda_1$ and $\Lambda$ with $m_t < \Lambda_1 < \Lambda$ by modifying eq. (12)

$$\Sigma_t(p^2) = \begin{cases} 
0 & \text{for } p^2 > \Lambda^2; \\
\sqrt{r} \cdot m_t & \text{for } \Lambda_1^2 \leq p^2 \leq \Lambda^2; \\
m_t & \text{for } p^2 < \Lambda_1^2,
\end{cases} \quad (15)$$
where $\Sigma$ is changed between $\Lambda_1$ and $\Lambda$. For $r > 1$ there is an extra “bump” between $\Lambda_1$ and $\Lambda$ which affects $\Delta \rho$. For $\Lambda^2, \Lambda_1^2 \gg m_t^2, rm_t^2$ we get

$$
\Delta \rho \simeq \frac{N_c m_t^2}{32\pi^2 v^2} \left( 1 + \frac{m_t^2}{\Lambda^2} - \left[ \frac{m_t^2 (\Lambda^2 - \Lambda_1^2)}{\Lambda^2 \Lambda_1^2} (r^2 - 1) \right] \right),
$$

(16)

where extra contributions due to $r \neq 1$ and $\Lambda_1 \neq \Lambda$ are isolated in square brackets. We can see that the bump counteracts the effect of the cutoff and makes $\Delta \rho$ less positive. In principle the bump can even be chosen to make $\Delta \rho$ vanish. The relation eq. (14) between $m_t$ and $M_W$ becomes also modified. For $\Lambda_2, \Lambda_2^2 \gg m_t^2, rm_t^2$ we get approximately

$$
M_W^2 \simeq \frac{g^2_N c}{32\pi^2} m_t^2 \left( \ln \left( \frac{\Lambda^2 - m_t^2}{m_t^2} \right) + \left[ (r - 1) \ln \left( \frac{\Lambda^2}{\Lambda_1^2} \right) \right] \right),
$$

(17)

where extra contributions due to the bump are again isolated in square brackets.

Now we add a slowly falling high energy tail to the last ansatz eq. (15)

$$
\Sigma_t(p^2) = \begin{cases} 
\text{equation (15)} & \text{for } p^2 < \Lambda^2; \\
\sqrt{r}m_t \left( \frac{p^2}{\Lambda^2} \right)^{\alpha} & \text{for } p^2 > \Lambda^2,
\end{cases}
$$

(18)

where $\alpha > 0$ is assumed. This high energy tail which is parametrized by $\alpha$ leads to

$$
\Delta \rho \simeq \frac{N_c m_t^2}{32\pi^2 v^2} \left( 1 + \frac{m_t^2}{\Lambda^2} - \left[ \frac{m_t^2 (\Lambda^2 - \Lambda_1^2)}{\Lambda^2 \Lambda_1^2} (r^2 - 1) \right] - \left\{ \frac{r^2}{4\alpha + 1} \frac{m_t^2}{\Lambda^2} \right\} \right),
$$

(19)

and

$$
M_W^2 = \frac{g^2_N c}{32\pi^2} m_t^2 \left( \ln \left( \frac{\Lambda^2 - m_t^2}{m_t^2} \right) + \left[ (r - 1) \ln \left( \frac{\Lambda^2}{\Lambda_1^2} \right) \right] + \left\{ \frac{r}{2\alpha} \right\} \right),
$$

(20)

where the extra corrections due to the tail are isolated in curly brackets.

Note that we are looking for a scenario which simultaneously avoids the fine-tuning problem and which is phenomenologically acceptable. Consequently $\Lambda$ and $\Lambda_1$ should be $TeV$-ish and the top mass values required from the $\Delta \rho$- and $M_W$-data should agree. This requires consequently some gap equation with a generic condensation scale $O(TeV)$ capable of producing a bump, and a tail – maybe of the type discussed in ref. [10]. The asymptotic high energy behaviour of $\Sigma_t$ might be described by a renormalization group equation if the spectrum of the theory does not contain further mass thresholds. This would imply a logarithmic tail and the parameter $\alpha$ should be very small. We could for example fix $\alpha$ in the minimal scenario by expanding the Higgs less one-loop renormalization group equation for $g_t$ in the Standard Model. This would lead to $\alpha \approx 0.04$. For such small values of $\alpha$ the tail leads to mild effects in the $\rho$-parameter and drastic changes in the $M_W$–$m_t$ relation.

We can illustrate the effects of the combined bump and tail by plotting eqs. (19) and (20) in Fig. 2 as solid lines for the parameters $r = 2$, $\Lambda = 2\Lambda_1$, $\Lambda_1 = 2m_t$ and $\alpha = 0.04$. The small value of $\alpha$ (corresponding to a logarithmic high energy tail of $\Sigma_t$) influences mostly...
the $M_W-m_t$ relation while the bump affects essentially only the $\Delta\rho-m_t$ relation. Taking into account experimental and theoretical errors the two top mass values agree for low values of $\Lambda$ consistent with the above assumptions and avoiding fine–tuning. We have thus illustrated that structured solutions of $\Sigma_t$ can solve the fine–tuning problem, i.e. allow for $\Lambda$–values within a few $\text{TeV}$. Furthermore the predicted $m_t-M_W-\Delta\rho$ relations are modified to be consistent with the data on $M_W$ and $\Delta\rho$. The predicted top mass differs however typically somewhat from its Standard Model value – something that will only be tested by a direct search for the top quark. A bump and a tail as discussed could for example be relevant in models of top condensation where heavy gauge bosons trigger condensation \cite{6} or in bootstrap scenarios where the $t$–channel effects of a composite Higgs are non–negligible \cite{15}.

There are other electro–weak observables which are sensitive to the top mass value like for example the $Z\bar{b}b$ vertex. If $m_t$ is replaced by $\Sigma_t$ in the relevant diagrams then one finds however that the top mass dependence is replaced by sensitivity to $\Sigma_t$ at low momenta. Thus in a first approximation these quantities depend essentially on the pole mass. There are however corrections which should become observable if high enough precision can be reached.

In summary we find that top condensation models are both phenomenological viable and natural if $\Sigma_t$ has suitable structure. The calculation of $\Sigma_t$ from first principles is however in general very difficult for a given model due to the non–perturbative nature of the relevant Schwinger–Dyson equation. But a reliable probe of the discussed effects will emerge when the Fermilab Collider starts to push the direct top mass limits into the Standard Model window. If the above ideas are relevant then the top quark should not be found precisely in the often cited Standard Model window but somewhat higher.
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Figures

\[ K \]

\[ \frac{i}{q^2} \] \[ \tilde{K} \]

Figure 1: The $W$ propagator in leading order $g_2^2$ and exact in the new non-perturbative interactions. Fermionic self-energies are represented as fat dots and the four-fermion Kernel $K$ is represented by a fat circle. In the second line the Kernel is split into Goldstone Boson contributions (which arise due to the broken global symmetries with some non-trivial vertex function) and $\tilde{K}$ (which has no further massless poles).
Figure 2: The predicted (pole) top mass $m_t$ versus $\Lambda$ using $\Delta \rho$ and $M_W$ as experimental input. The upper dashed line follows from eq. (14) and the lower dashed line from eq. (13). The solid lines follow from the combined bump and tail ansatz for $\Sigma_t$ showing that low values of $\Lambda$ are then possible.