Particle Physics Colloquium KIT Karlsruhe — 13 December, 2018

& NCTS Annual Theory Meeting, Taiwan — 19 December, 2018

[extended version]

High-energy Higgsplosion

Valya Khoze

IPPP Durham

- VVK & Spannowsky 1704.03447, 1707.01531, 1809.11141
- VVK 1806.05648 & 1705.04365
- VVK & Reiness: Semiclassical Review 1810.01722

Most recent

In the first part of this talk: I'll outline the main idea of Higgsplosion

- Consider n~150 Higgs bosons produced in a final state at n lambda >> 1. Kinematically possible for scattering at E ~100 TeV

- HIGGSPLOSION: n-particle rates computed in a weakly-coupled theory become unsuppressed above certain critical values of n and $E = \sqrt{s} = \sqrt{p^2}$

$$m\,\Gamma_n(p^2) \;=\; \mathrm{Im}_n\Sigma(p^2) \;>\; p^2\,,\, m^2$$

In the second part of this talk:

- we'll consider an intrinsically Non-perturbative — semiclassical set-up $n \,\propto\, \sqrt{s}/m \,\propto\, 1/\lambda \,\gg\, 1$

- it incorporates correctly the tree-level results and

- the leading-order quantum effects = leading loops

already known

- compute quantum effects in the large lambda n limit

new

Conclusions & summary

Higgsplosion: *few* particles —> *many* particles processes

$$\sqrt{s}\ :\quad \underset{\textit{few}}{X} \;\to\; \underset{\textit{many}}{n \times \phi}\,,$$

$$\text{Scattering process :}\qquad |X(\sqrt{s})\rangle = |2\rangle \;\to\; |n\rangle \quad \Rightarrow \quad \text{cross section } \sigma_n(\sqrt{s})\,,$$

$$\text{Resonance decay :}\qquad \boxed{|X(\sqrt{s})\rangle = |1^*\rangle \;\to\; |n\rangle} \quad \Rightarrow \quad \text{partial width } \Gamma_n(s)\,.$$

A non-perturbative semiclassical approach can be used to compute such processes. The semiclassical approach assumes that the initial state X can be approximated by a point-like injection of energy: via a local operator O(x)

$$|X\rangle \;=\; \mathcal{O}(0)\,|0\rangle$$

Ideal for 1* —> n   Higgsplosion is when: $\frac{m}{s}\,\Gamma_n(s) \gtrsim 1$

For for 2 —> n OK for s-channel but not for t-channel where there is an impact parameter

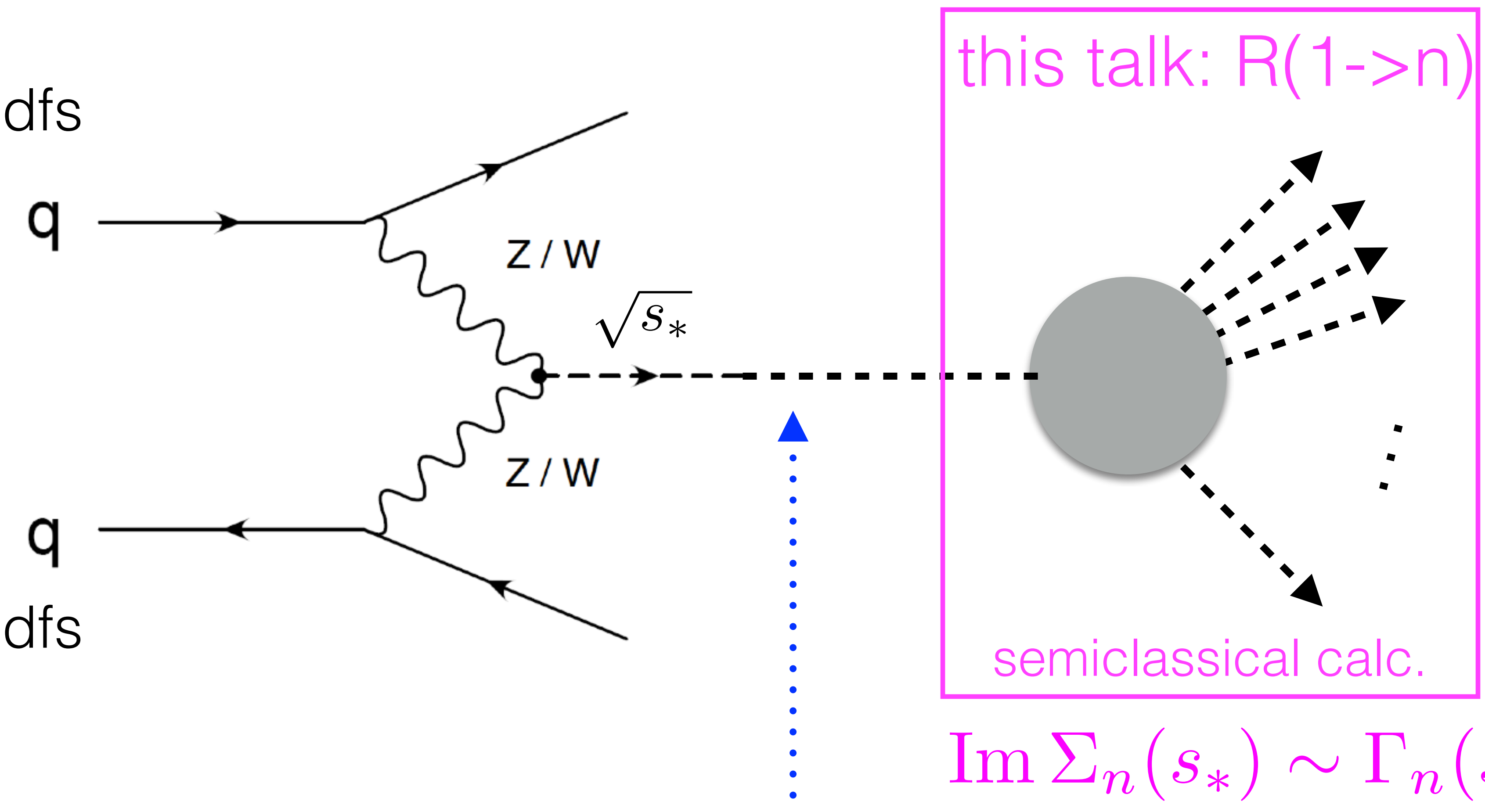
1)

1->n processes of interest for Higgsplosion

e.g.:Vector boson fusion in high-energy pp collisions at ~100 TeV

quark pdfs

$\sqrt{s'}$



quark pdfs

n non-relativistic Higgses Higgsplosion at $\sqrt{s_*}$

$$\mathrm{Im}\,\Sigma_n(s_*) \sim \Gamma_n(s_*)$$

$$\frac{i}{s_* - m_h^2 - Re\tilde{\Sigma}(s_*) + i m_h \Gamma(s_*)}$$

Propagator with Higgspersion at $\sqrt{s_*}$

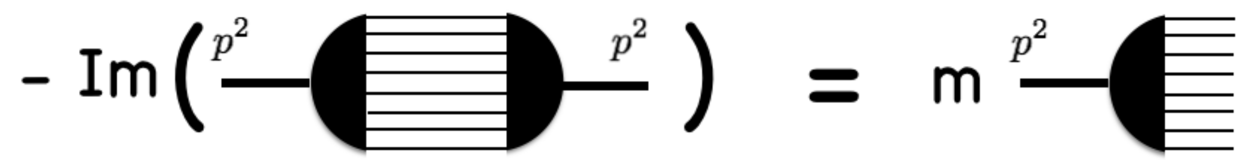
- VVK & Spannowsky 1704.03447,1707.01531

2)

Schwinger-propagator and optical theorem

The optical theorem relates the 1* -> n h amplitudes to the imaginary part of the self-energy (valid to all orders)

$$-\,\mathrm{Im}\,\Sigma_R(p^2) \;=\; m\,\Gamma(p^2)$$

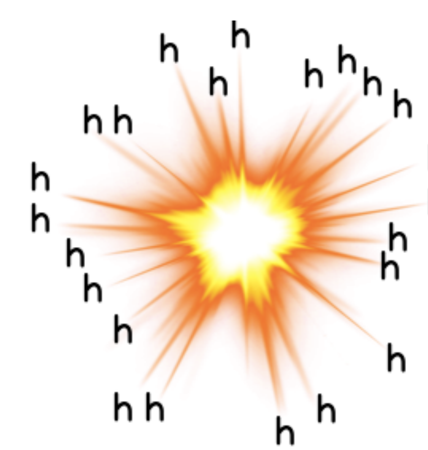


where $\Gamma(s) \;=\; \sum_{n=2}^{\infty} \Gamma_n(s)$ and $\Gamma_n(s) \;=\; \frac{1}{2m} \int \frac{d\Phi_n}{n!}\, |\mathcal{M}(1 \to n)|^2$

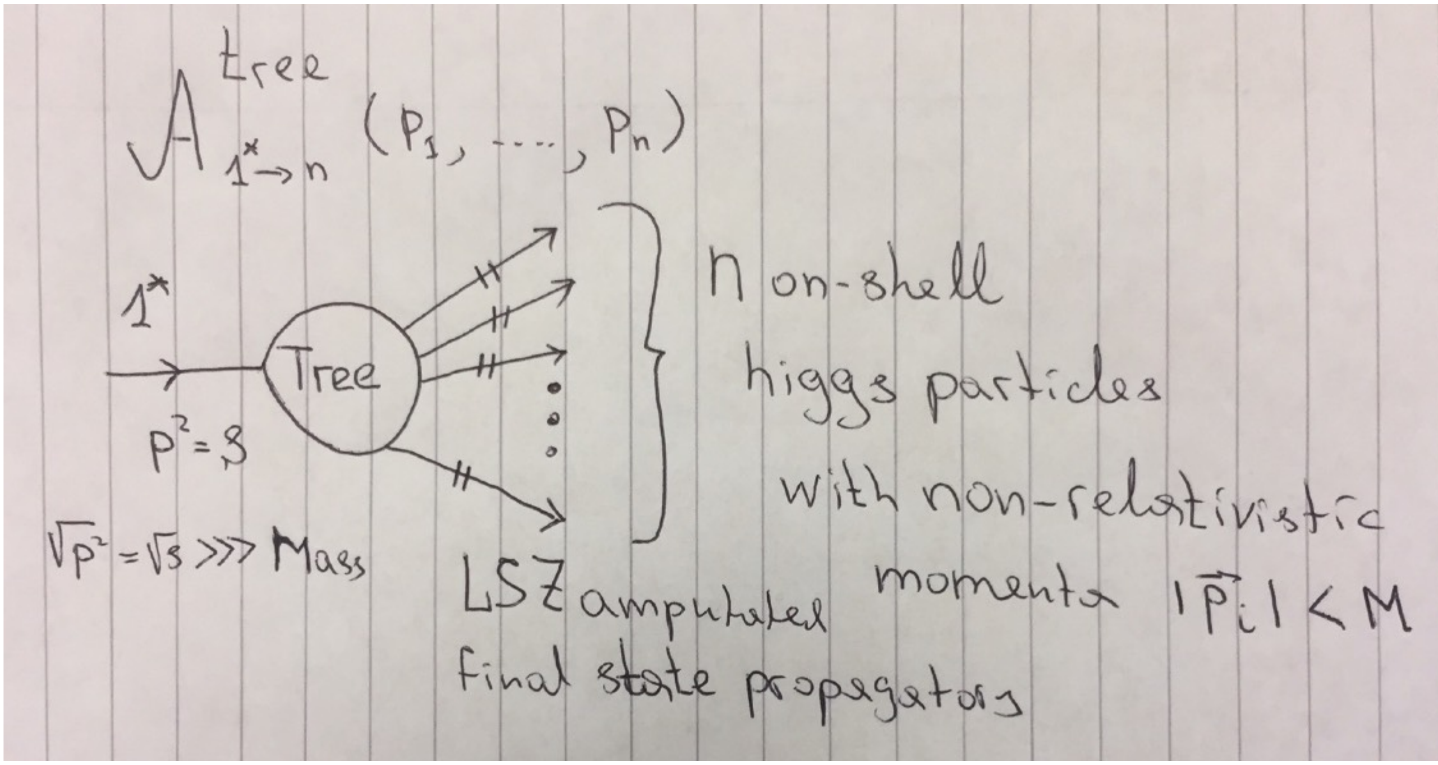
and thus $$\Delta_R(p) \;=\; \frac{i}{p^2 - m^2 \;-\; \mathrm{Re}\,\Sigma_R(p^2) \;+\; im\Gamma(p^2) \;+ i\epsilon}$$

No non-pert. information
about the Real part of Sigma
but it cannot cancel
Imaginary part

Higgsplodes
when
$m\Gamma_n(p^2) \gtrsim p^2$



Warm-up: Compute 1 -> n amplitudes@LO with non-relativistic final momenta:



see classic 1992-1994 papers:
Brown; Voloshin;
Argyres, Kleiss, Papodopoulos
Libanov, Rubakov, Son, Troitski

more recently: VVK 1411.2925

$$\mathcal{L} = \frac{1}{2}(\partial_\mu h)^2 - \frac{\lambda}{4}(h^2 - v^2)^2$$

prototype of the SM Higgs
in the unitary gauge

Tree-level $1^* \to n$ amplitudes in the limit $\varepsilon \to 0$ for any n are given by

$$\mathcal{A}_n(p_1, \ldots p_n) = \boxed{n! \left(\frac{\lambda}{2M_h^2}\right)^{\frac{n-1}{2}}} \left(1 - \frac{7}{6}n\varepsilon - \frac{1}{6}\frac{n}{n-1}\varepsilon + \mathcal{O}(\varepsilon^2)\right)$$

factorial growth

amplitude on the n-particle threshold

$$\varepsilon = \frac{1}{n\,M_h} E_n^{\rm kin} = \frac{1}{n}\frac{1}{2M_h^2}\sum_{i=1}^n \vec{p}_i^{\,2}$$

kinetic energy per particle per mass

In the large-n-non-relativistic limit the result is

$$\boxed{\mathcal{A}_n(p_1, \ldots p_n) = n! \left(\frac{\lambda}{2M_h^2}\right)^{\frac{n-1}{2}} \exp\left[-\frac{7}{6}n\varepsilon\right], \quad n \to \infty,\ \varepsilon \to 0,\ n\varepsilon = \text{fixed}}$$

Square the amplitude & integrate over the n-particle phase-space:

The cross-section and/or the n-particle partial decay Γ_n

$$\Gamma_n(s) \;=\; \int d\Phi_n \, \frac{1}{n!} \, \left|\mathcal{A}_{h^* \to n \times h}\right|^2$$

The n-particle Lorentz-invariant phase space volume element

$$\int d\Phi_n \;=\; (2\pi)^4 \delta^{(4)}(P_{\rm in} - \sum_{j=1}^{n} p_j) \, \prod_{j=1}^{n} \int \frac{d^3 p_j}{(2\pi)^3 \, 2p_j^0} \, ,$$

in the large-n non-relativistic limit with $n\varepsilon_h$ fixed becomes,

$$\Phi_n \;\simeq\; \frac{1}{\sqrt{n}} \left(\frac{M_h^2}{2}\right)^n \; \exp\left[\frac{3n}{2}\left(\log\frac{\varepsilon_h}{3\pi} + 1\right) \;+\; \frac{n\varepsilon_h}{4} \;+\; \mathcal{O}(n\varepsilon_h^2)\right]$$

We find:

$$\boxed{\Gamma_n^{\rm tree}(s) \;\sim\; \exp\left[n\left(\log\frac{\lambda n}{4} - 1\right) \;+\; \frac{3n}{2}\left(\log\frac{\varepsilon}{3\pi} + 1\right) \;-\; \frac{25}{12}\, n\varepsilon \;+\; \mathcal{O}(n\varepsilon^2)\right]}$$

Son 1994;
Libanov, Rubakov, Troitskii 1997; more recently: VVK 1411.2925

Can also include *loop corrections* to amplitudes on thresholds:

The 1-loop corrected threshold amplitude for the pure n Higgs production:

$$\boxed{\phi^4 \text{ with SSB}: \quad \mathcal{A}_{1\to n}^{\text{tree+1loop}} = n!\,(2v)^{1-n}\left(1 + n(n-1)\frac{\sqrt{3}\lambda}{8\pi}\right)}$$

There are strong indications, based on the analysis of leading singularities of the multi-loop expansion around singular generating functions in scalar field theory, that the 1-loop correction exponentiates,

Libanov, Rubakov, Son, Troitsky 1994

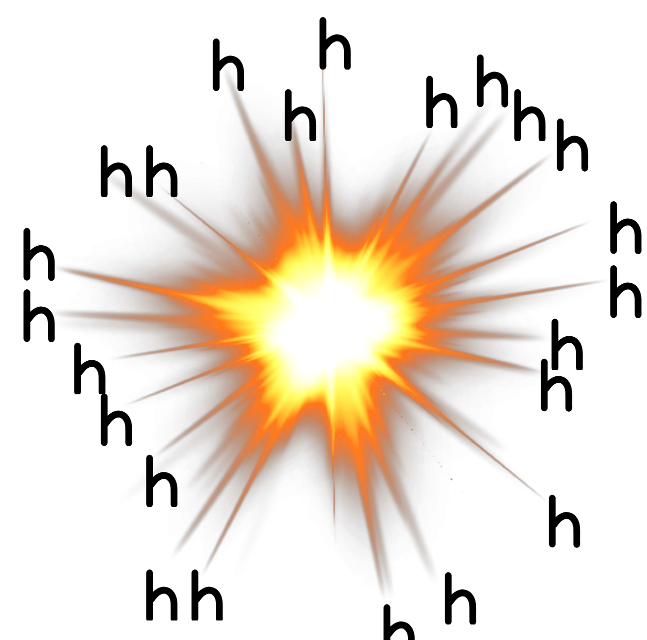
$$\boxed{\mathcal{A}_{1\to n} = \mathcal{A}_{1\to n}^{\text{tree}} \times \exp\left[B\,\lambda n^2 + \mathcal{O}(\lambda n)\right]}$$

in the limit $\lambda \to 0$, $n \to \infty$ with λn fixed. Here B is determined from the 1-loop calculation (as above) – *Smith; Voloshin 1992*):

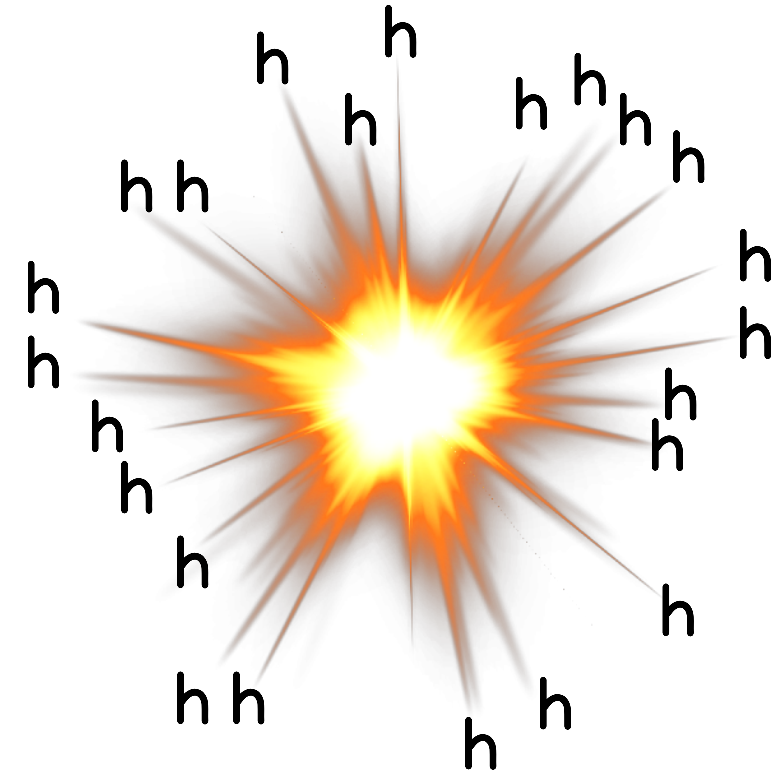
$$B = +\frac{\sqrt{3}}{4\pi}$$

Really need to switch to the regime of lambda n >>1

For this we need a non-perturbative —
semiclassical approach — next slide…



Explosive growth of 1->n perturbative process



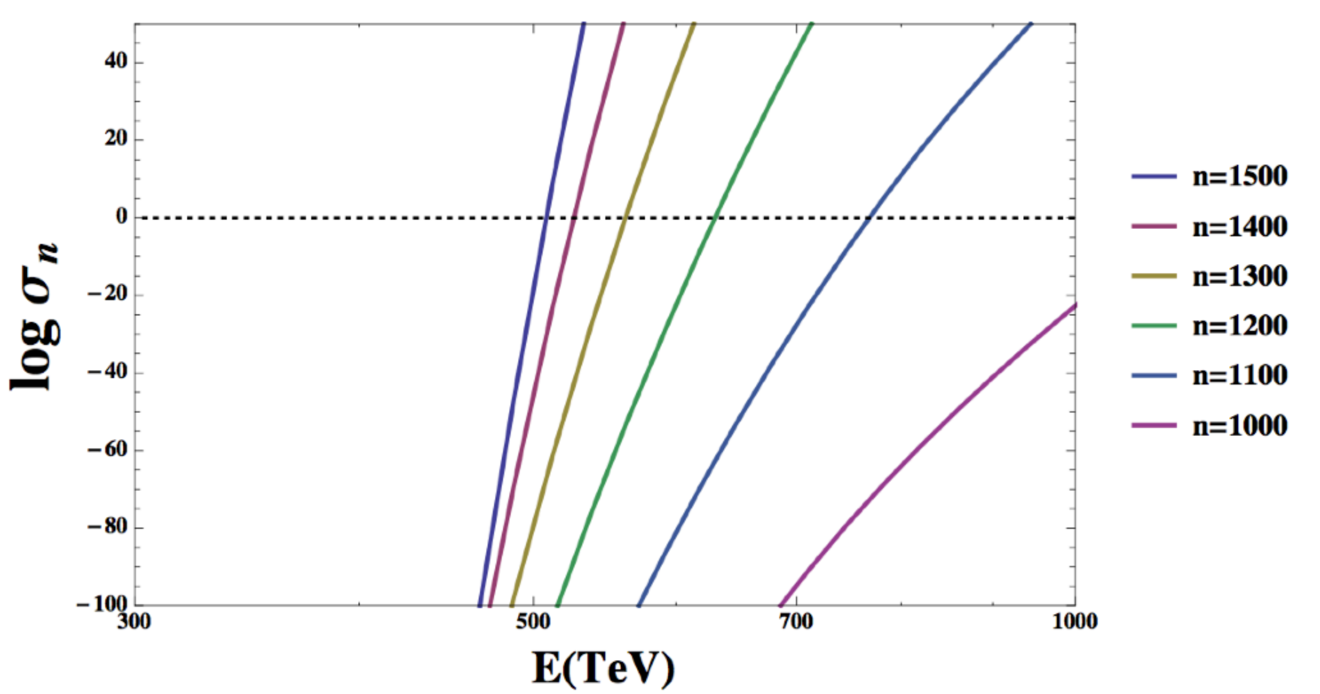
$$\mathcal{R}_n(s) \;:=\; \frac{1}{2M_h^2}\int d\Pi_n |\mathcal{M}(1\to n)|^2$$

$$\mathcal{R}(\lambda; n, \varepsilon) \;=\; \exp\left[n\left(\log\frac{\lambda n}{4} - 1\right) + \frac{3n}{2}\left(\log\frac{\varepsilon}{3\pi}+1\right) - \frac{25}{12}n\varepsilon\right]$$

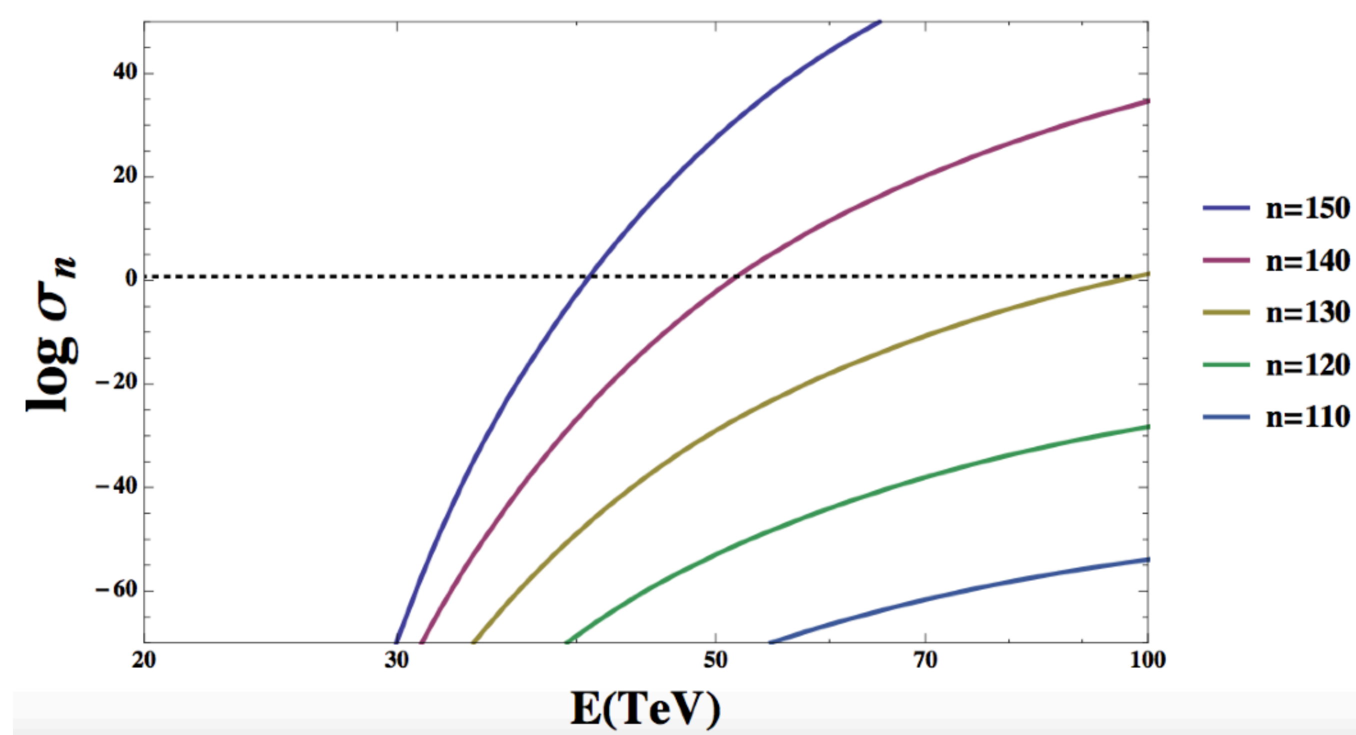
Extreme energy dependence for 1^* -> n cross section

including 1-loop result reduces 'ignition' scale

$\log \sigma_n^{\text{tree}}$



$\log \sigma_n^{\text{loop}}$

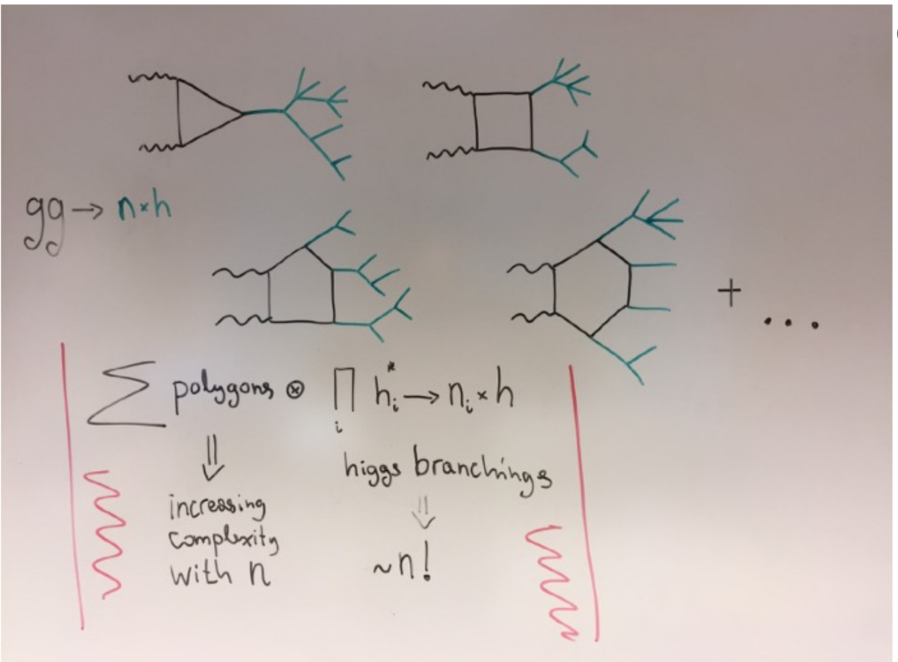


Higgspersion of the propagator due to Im Sigma not yet included here!

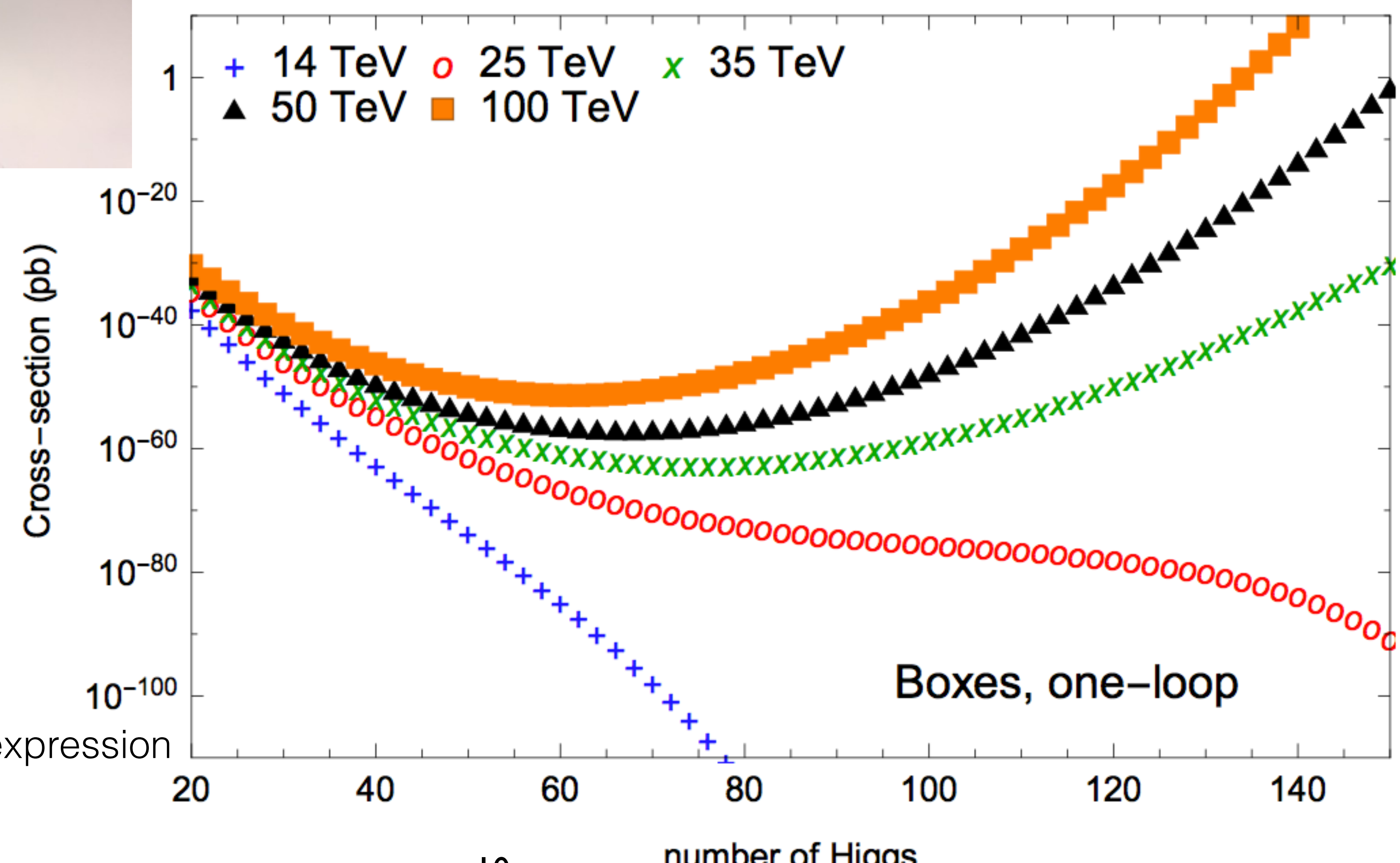
We'll need a non-perturbative semiclassical approach: Part 2 of the talk

Application to gluon fusion:

It was argued that these results can be used to assess what collider energy needed to test where perturbation theory becomes strong [in gluon fusion]



Degrande, VVK, Mattelaer '16



using the 1-loop improved expression

Factorial growth of tree-level amplitudes at threshold is captured by classical solutions

$$\mathcal{L}(h) \;=\; \frac{1}{2}\,(\partial h)^2 - \frac{\lambda}{4}\left(h^2 - v^2\right)^2 \;,$$

The classical equation for the spatially uniform field $h(t)$,

$$d_t^2 h \;=\; -\lambda\, h^3 + \lambda v^2\, h \,,$$

has a closed-form solution with correct initial conditions $h_{\rm cl} = v + z + \ldots$

$$h_0(z_0;t) \;=\; v\,\left(\frac{1 + z_0\, e^{imt}/(2v)}{1 - z_0\, e^{imt}/(2v)}\right)\,, \quad m = \sqrt{2\lambda} v$$

$$h_0(z) \;=\; v \,+\, 2v \sum_{n=1}^{\infty} \left(\frac{z}{2v}\right)^n \,, \quad z = z(t) = z_0\, e^{imt}$$

$$\mathcal{A}_{1\to n} \;=\; \left(\frac{\partial}{\partial z}\right)^n h_{\rm cl}\bigg|_{z=0} = \boxed{n!\,(2v)^{1-n}}$$

Factorial growth

L. Brown 9209203

Classical Solutions & singularities in complex time:

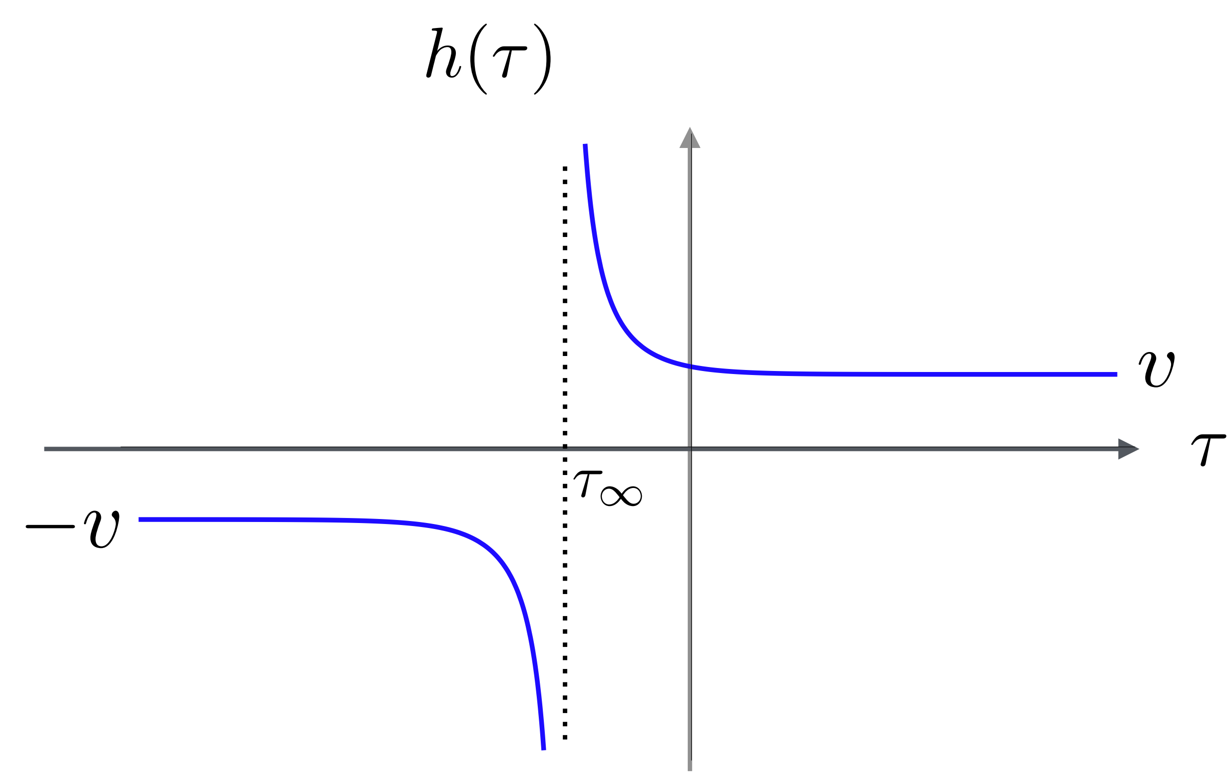
$$t \longrightarrow t_{\mathbb{C}} = t + i\tau$$

$$h_0(t_{\mathbb{C}}) = v \left(\frac{1 + e^{im(t_{\mathbb{C}} - i\tau_\infty)}}{1 - e^{im(t_{\mathbb{C}} - i\tau_\infty)}} \right),$$

$$\tau_\infty := \frac{1}{m} \log \left(\frac{z_0}{2v} \right)$$

Our simple example of a classical solution
(corresponding to the tree-level Amplitudes)

$$h_0(\tau) = v \left(\frac{1 + e^{-m(\tau - \tau_\infty)}}{1 - e^{-m(\tau - \tau_\infty)}} \right)$$



Such singular complex-valued solutions will emerge in the semiclassical approach

Part II

Main idea of the semiclassical approach

$\mathcal{R}_n(E)$ is the probability rate for a local operator $\mathcal{O}(0)$ to create n particles of total energy E from the vacuum,

$$\mathcal{R}_n(E) \; = \; \int \frac{1}{n!} d\Phi_n \, \langle 0| \, \mathcal{O}^\dagger \, S^\dagger \, P_E |n\rangle\langle n| \, P_E \, S \, \mathcal{O} \, |0\rangle$$

P_E is the projection operator on states with fixed energy E.

$$\mathcal{O} \; = \; e^{j\, h(0)} \, ,$$

and the limit $j \to 0$ is taken in the computation of the probability rates,

$$\mathcal{R}_n(E) \; = \; \lim_{j\to 0} \int \frac{1}{n!} d\Phi_n \, \langle 0| \, e^{j\, h(0)}{}^\dagger \, S^\dagger \, P_E |n\rangle\langle n| \, P_E \, S \, e^{j\, h(0)} \, |0\rangle \, .$$

Note: non-dynamical (non-propagating) initial state $\mathcal{O}|0\rangle$.
The semi-classical (steepest descent) limit:

$$\varepsilon = \frac{E - nm}{nm}$$

$$\boxed{\lambda \to 0 \, , \quad n \to \infty \, , \quad \text{with} \quad \lambda n = \text{fixed} \, , \quad \varepsilon = \text{fixed} \, .}$$

Evaluate the path integral in this double-scaling limit.
n enters via the coherent state formalism.

Rubakov & Tinyakov; DT Son '95

Main idea of the semiclassical approach

Note (1):

The initial state is not semiclassical, it contains few rather than many particles.

Rubakov et al & Son argued that it can be approximated in the semiclassical method by a certain local operator acting on the vacuum:

$$|X\rangle = \mathcal{O}(0)\,|0\rangle$$

$$\mathcal{O}(x) = j^{-1}\, e^{j\phi(x)}\,,$$

j is a constant $j = c/\lambda$. Finally one takes the limit $c \to 0$ (or equivalently $j \to 0$)

A refinement:

smear O(x) with a wave packet / test function

operator localized in the vicinity of a point x

$$\mathcal{O}_g(x) = \int d^4x'\, g(x'-x)\,\mathcal{O}(x')\,,$$

$$|X\rangle = \mathcal{O}_g(0)\,|0\rangle = \int d^4x'\, g(x')\,\mathcal{O}(x')\,|0\rangle$$

The Semiclassical formalism of Son: results in four steps

1. Solve the classical equation without the source-term:

$$\frac{\delta S}{\delta h(x)} = 0$$

   a complex-valued solution $h(x)$ with a point-like singularity at $x^\mu = 0$. The singularity is due to $\mathcal{O}(x=0)$.

2. Impose the initial and final-time boundary conditions:

$$\lim_{t\to-\infty} h(x) = v + \int \frac{d^3k}{(2\pi)^{3/2}} \frac{1}{\sqrt{2\omega_{\mathbf{k}}}} \, a^\dagger_{\mathbf{k}} \, e^{ik_\mu x^\mu}$$

$$\lim_{t\to+\infty} h(x) = v + \int \frac{d^3k}{(2\pi)^{3/2}} \frac{1}{\sqrt{2\omega_{\mathbf{k}}}} \left(b_{\mathbf{k}} \, e^{\omega_{\mathbf{k}} T - \theta} \, e^{-ik_\mu x^\mu} + b^\dagger_{\mathbf{k}} \, e^{ik_\mu x^\mu} \right)$$

- Son hep-ph/055338

The Semiclassical formalism of Son: results in four steps

3. Compute E and n of the final state using the $t \to +\infty$ asymptotics

$$E = \int d^3k\ \omega_{\mathbf{k}}\, b^\dagger_{\mathbf{k}}\, b_{\mathbf{k}}\, e^{\omega_{\mathbf{k}} T - \theta}\,, \qquad n = \int d^3k\ b^\dagger_{\mathbf{k}}\, b_{\mathbf{k}}\, e^{\omega_{\mathbf{k}} T - \theta}$$

At $t \to -\infty$ the energy and the particle number are vanishing.
The energy changes discontinuously from 0 to E at the singularity at $t = 0$.

4. Eliminate the T and θ parameters in favour of E and n.
Finally, compute the function $W(E,n)$

$$\boxed{W(E,n) = ET - n\theta - 2\mathrm{Im}S[h]}$$

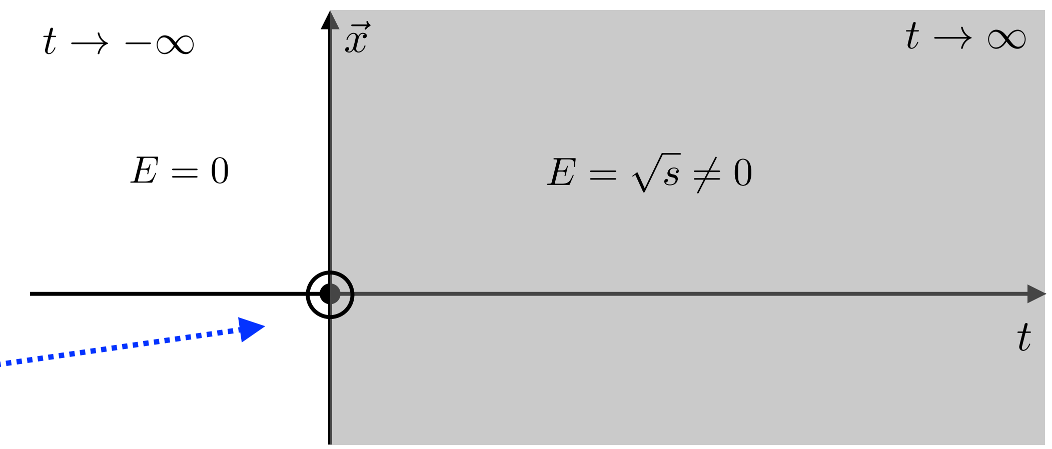
on the set $\{h(x), T, \theta\}$ and fine the semiclassical rate $\mathcal{R}_n(E) = \exp\left[W(E,n)\right]$

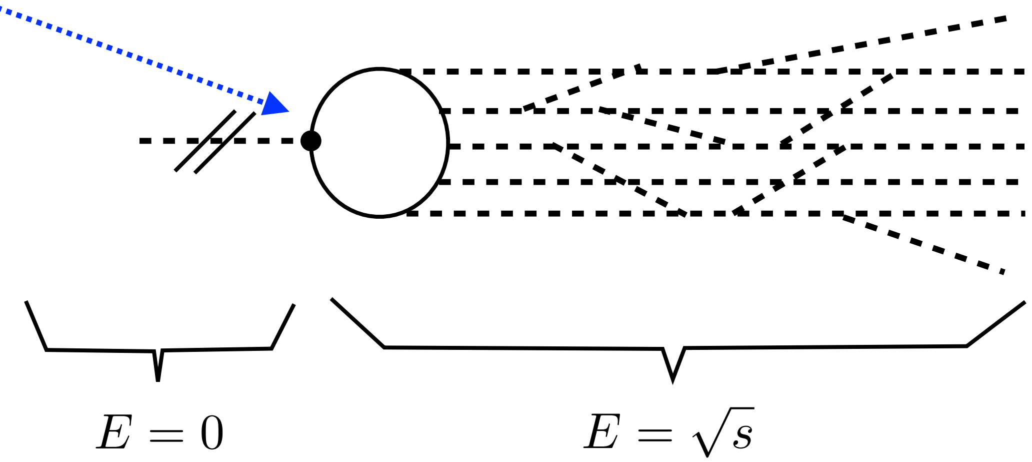
- Son hep-ph/055338

Main idea of the semiclassical approach

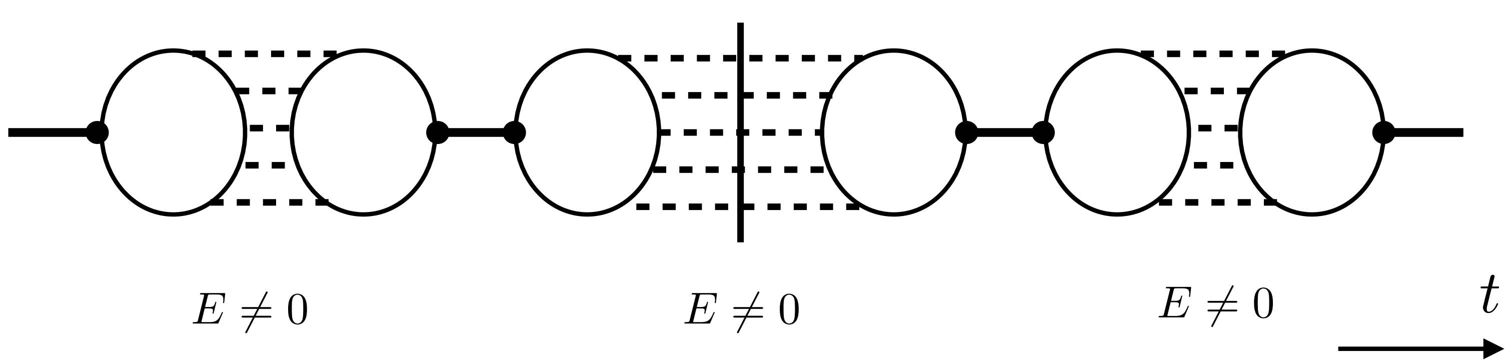
Note (2): $\mathrm{Im}\,\Sigma(s)$

The classical solutions that we use have a single point-like singularity in Minkowski space at the point x=0 where the operator O(0) is located.





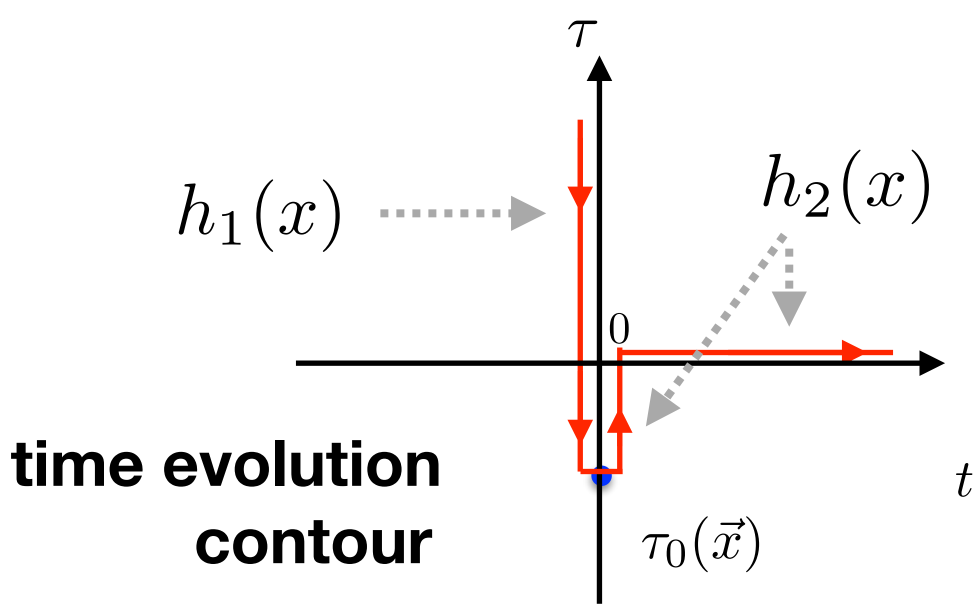
Such configurations contribute to 1PI matrix elements i.e. precisely to $\mathrm{Im}\,\Sigma(s)$

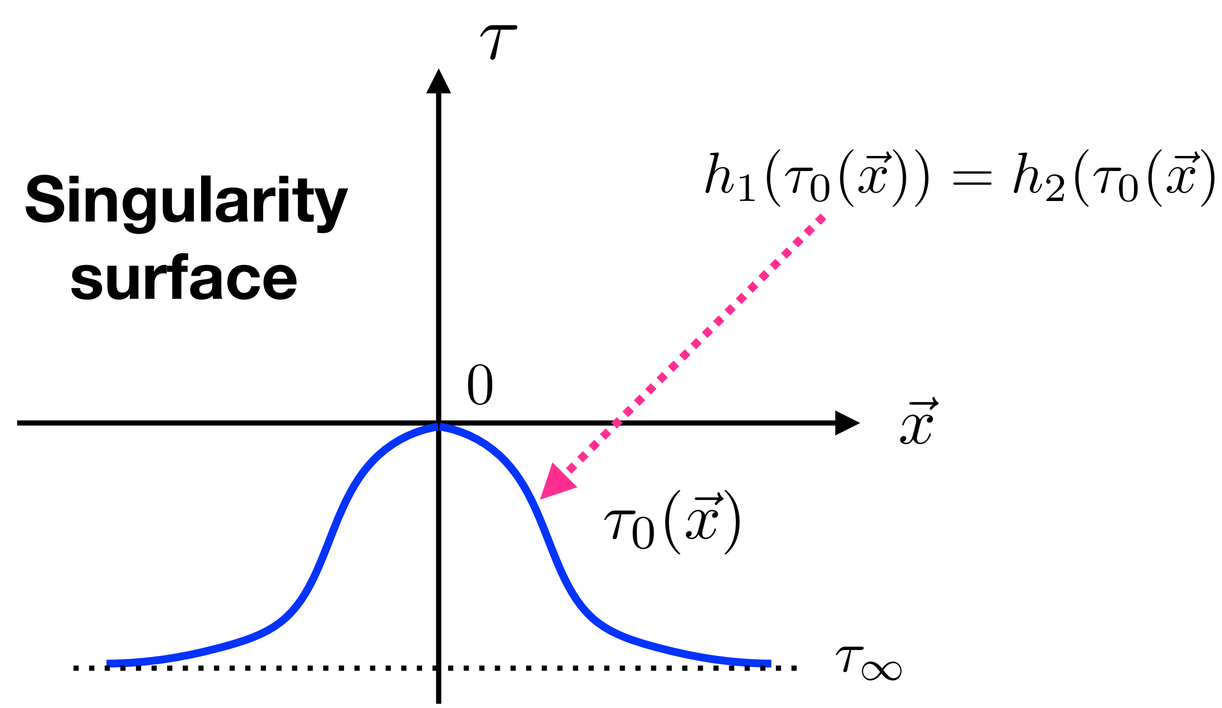


1-particle-reducible contributions to would require multiple singularities, i.e. multiple energy jumps.

Refining the method in complex time

- In Euclidean space-time $(\tau, \vec{x})$ the solution is singular on a 3-dimensional surface $\tau = \tau_0(\vec{x})$.





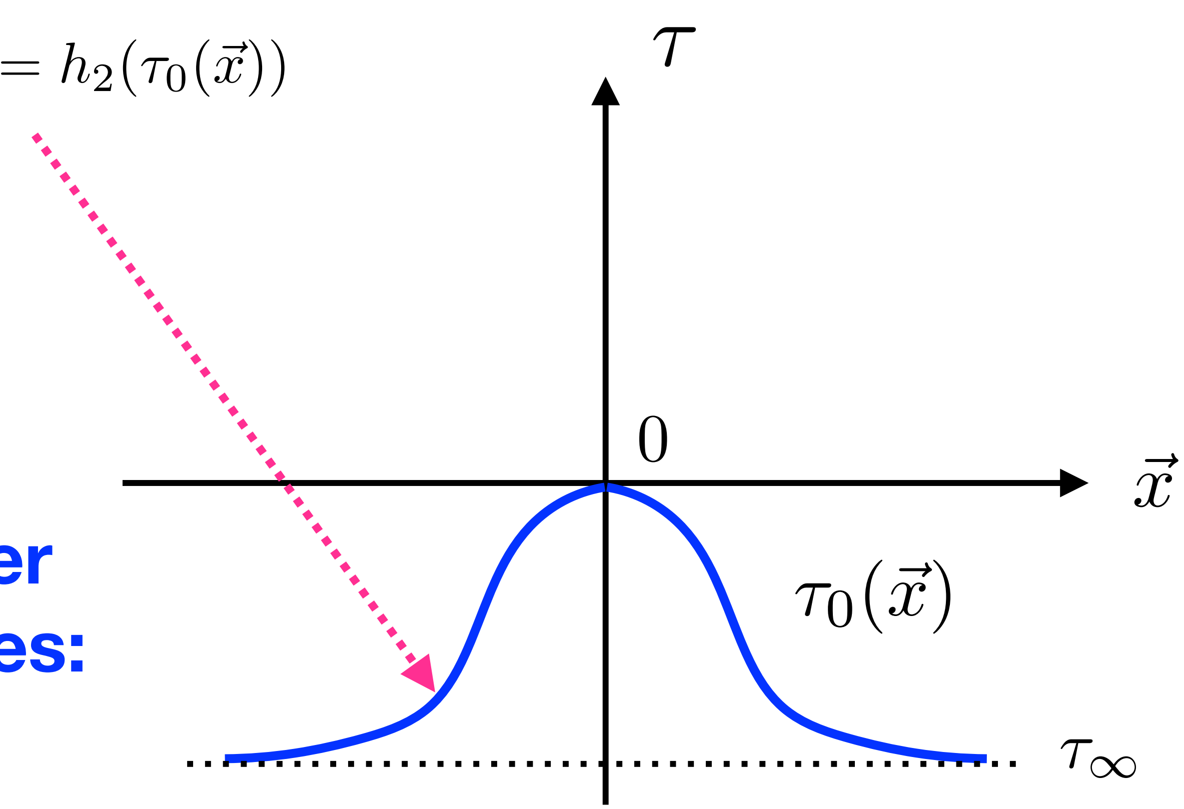
- Find a classical trajectory $h_1(\tau, \vec{x})$ satisfying initial time boundary cond-s.
- Find another classical trajectory $h_2(\tau, \vec{x})$ satisfying final time conditions.
- h_1 and h_2 are singular on $\tau_0(\vec{x})$ and $h_1(\tau_0(\vec{x}), \vec{x}) = h_2(\tau_0(\vec{x}), \vec{x})$
- Extremize the action S over all singularity surfaces $\tau_0(\vec{x})$.

- For the combined configuration $h(x)$ to solve classical equations everywhere, including the τ_0 surface:

need to extremize the action integral over all singularity surfaces $\tau = \tau_0(\vec{x})$ containing the point $t = 0 = \vec{x}$.

$$iS[h] = \int d^3x \left(\left| \int_{+\infty}^{\tau_0(\vec{x})} d\tau \, \mathcal{L}_{\mathrm{Eucl}}(h_1) \right| - \left| \int_{\tau_0(\vec{x})}^{0} d\tau \, \mathcal{L}_{\mathrm{Eucl}}(h_2) \right| + i \int_0^\infty dt \, \mathcal{L}(h_2) \right)$$

$h_1(\tau_0(\vec{x})) = h_2(\tau_0(\vec{x}))$



Extremize the action S over all such singularity surfaces:

Computing the semiclassical rate

Classical solution singular on generic tau_0(x) surfaces:

$$h(t_{\mathbb{C}},\vec{x}) \;=\; v\left(\frac{1\;+\;e^{im(t_{\mathbb{C}}-i\tau_\infty)}}{1\;-\;e^{im(t_{\mathbb{C}}-i\tau_\infty)}}\right)\;+\;\tilde{\phi}(t_{\mathbb{C}},\vec{x})$$

Find that:

$$\boxed{\begin{aligned} W(E,n) \;&=\; ET \;-\; n\theta \;-\; 2\mathrm{Re}S_{\mathrm{Eucl}}[h] \\ &=\; n\,\log\frac{\lambda n}{4}\;+\;\frac{3n}{2}\left(\log\frac{3\pi}{\varepsilon}+1\right)\;-\;2nm\,\tau_\infty\;-\;2\mathrm{Re}S_{\mathrm{Eucl}}[h] \end{aligned}}$$

$W(E,n)^{\mathrm{tree}}$

agrees with the known result
of tree-level contributions

$\Delta W^{\mathrm{quant}}$

need to compute by extremizing
w.r.t tau_0(x)

Computing the semiclassical rate

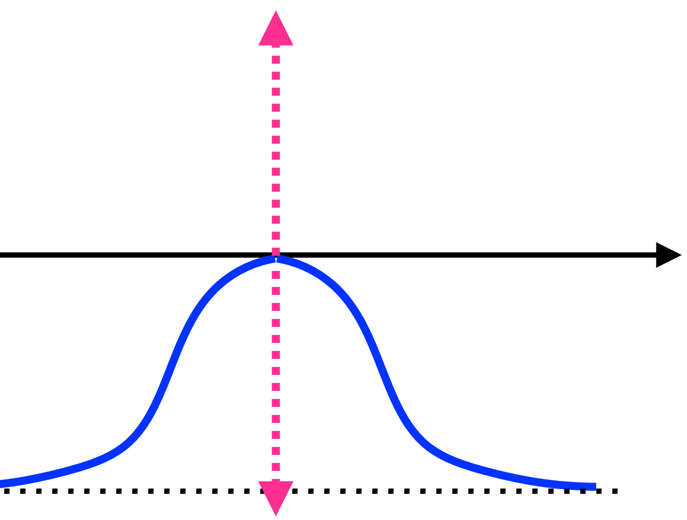
$$\Delta W^{\rm quant} = -2nm\,\tau_\infty - 2{\rm Re}\,S^{(1,2)}_{\rm Eucl}$$
$$= 2nm\,|\tau_\infty| + 2\int d^3x \left[\int_{\tau_0(\vec{x})}^{+\infty} d\tau\,\mathcal{L}_{\rm Eucl}(h_1) - \int_{\tau_0(\vec{x})}^{0} d\tau\,\mathcal{L}_{\rm Eucl}(h_2)\right]$$

Use *thin wall* approximation:

Force x height

E=0 configuration

E=mn configuration



$$\frac{1}{2}\Delta W^{\rm quant} = nm\,|\tau_\infty| - \underbrace{\int_{A+i\epsilon}^{0+i\epsilon} d\tau\,L_{\rm Eucl}(h_2;\tau_0(\vec{x}))}_{\equiv S_{\rm Eucl}[\tau_0(\vec{x})]} + \frac{4\pi}{3}\mu R^3$$

Surface-energy

Force x height

Surface-energy

Mechanical analogy: surface at equilibrium of forces.

- Gorsky & Voloshin hep-ph/9305219
- VVK 1806.05648

Computing the semiclassical rate for $\lambda n \gg 1$

Use *thin wall* approximation:

$$S_{\rm Eucl}[\tau_0(r)] \;=\; \int_{\tau_\infty}^{0} d\tau\, 4\pi\mu\, r^2\sqrt{1+\dot r^2} \;\equiv\; \int_{\tau_\infty}^{0} d\tau\, L(r,\dot r)$$

Surface tension

$$\mu \;=\; \int_{-\infty-i\epsilon}^{+\infty-i\epsilon} d\tau \left(\frac{1}{2}\left(\frac{dh}{d\tau}\right)^2 + \frac{\lambda}{4}\left(h^2-v^2\right)^2\right) \;=\; \frac{m^3}{3\lambda}$$

Conjugate momentum

$$p \;=\; \frac{\partial L(r,\dot r)}{\partial \dot r} \;=\; 4\pi\,\mu\frac{r^2\dot r}{\sqrt{1+\dot r^2}}$$

Hamiltonian => Energy

$$H(p,r) \;=\; L(r,\dot r) \;-\; p\,\dot r$$

$$\frac{1}{2}\Delta W^{\rm quant} \;=\; (E-nm)\tau_\infty \;-\; \int_R^0 p(E)\,dr \;+\; \frac{4\pi}{3}\,\mu R^3$$

Quantum rate on the stationary trajectory:

$$\frac{1}{2}\Delta W^{\rm quant}_{\rm stationary} \;=\; -\int_R^0 p(E)\,dr \;+\; \frac{4\pi}{3}\,\mu R^3\,, \qquad E = nm$$

- Gorsky & Voloshin hep-ph/9305219
- VVK 1806.05648

Computing the semiclassical rate

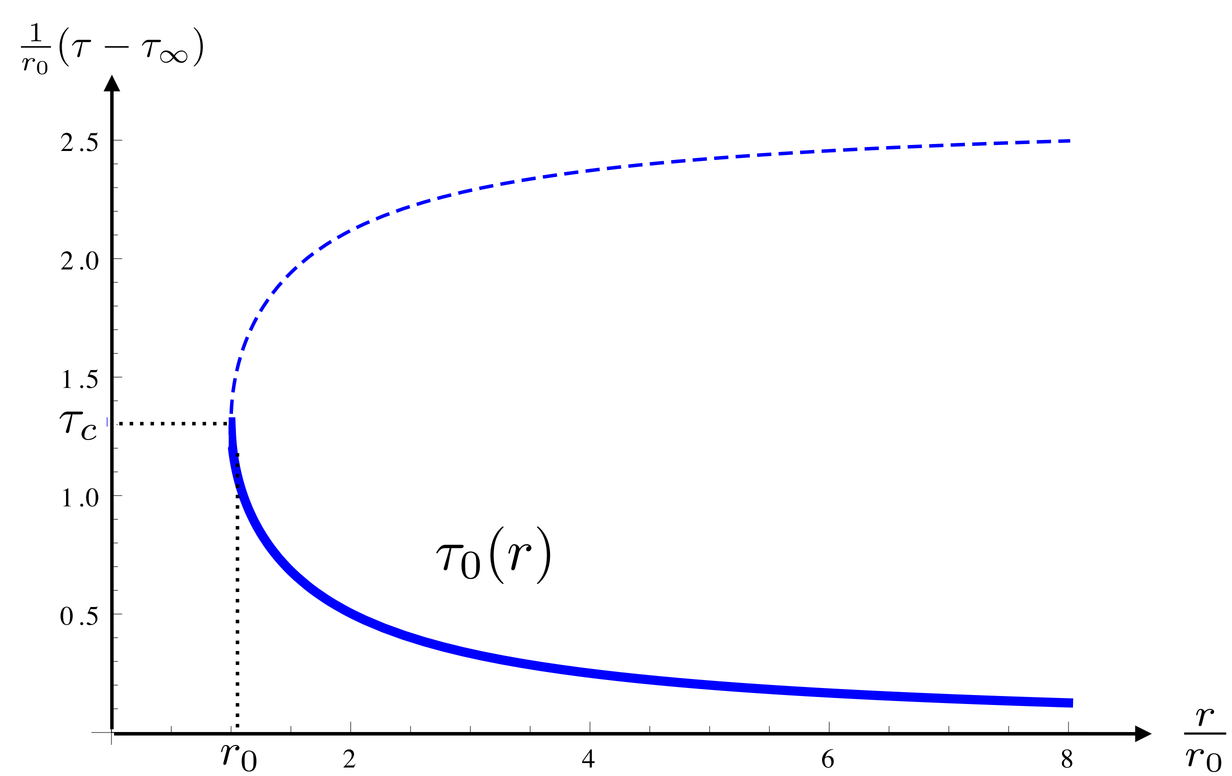
Use *thin wall* approximation:

$$\frac{1}{2}\Delta W^{\rm quant}_{\rm stationary} = -\int_R^0 p(E)\,dr + \frac{4\pi}{3}\,\mu R^3\,, \qquad E = nm$$

final result

$$\boxed{\Delta W^{\rm quant} = \frac{E^{3/2}}{\sqrt{\mu}}\,\frac{2}{3}\,\frac{\Gamma(5/4)}{\Gamma(3/4)} = \frac{1}{\lambda}\,(\lambda n)^{3/2}\,\frac{2}{\sqrt{3}}\,\frac{\Gamma(5/4)}{\Gamma(3/4)} \simeq 0.854\,n\sqrt{\lambda n}}$$

Classical trajectory tau(r):



Justifies the thin wall approximation:

$$rm \geq r_0 m = m\left(\frac{E}{4\pi\mu}\right)^{1/2} \propto \left(\frac{\lambda\,E}{m}\right)^{1/2} = \sqrt{\lambda n} \gg 1\,, \qquad r_0^2 = \frac{E}{4\pi\mu}$$

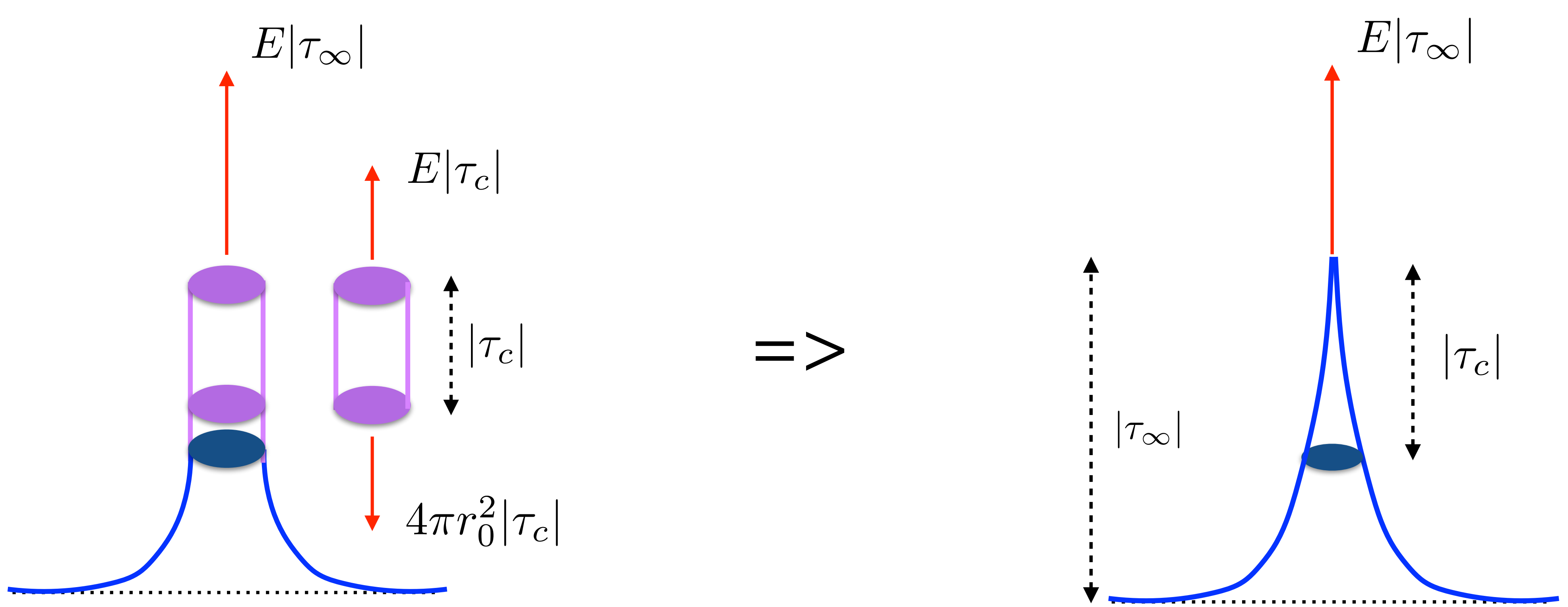
Computing the semiclassical rate

Use *thin wall* approximation:

$$\frac{1}{2}\Delta W^{\text{quant}}_{\text{stationary}} = -\int_R^0 p(E)\,dr + \frac{4\pi}{3}\,\mu R^3\,, \qquad E = nm$$

final result

$$\boxed{\Delta W^{\text{quant}} = \frac{E^{3/2}}{\sqrt{\mu}}\,\frac{2}{3}\,\frac{\Gamma(5/4)}{\Gamma(3/4)} = \frac{1}{\lambda}\,(\lambda n)^{3/2}\,\frac{2}{\sqrt{3}}\,\frac{\Gamma(5/4)}{\Gamma(3/4)} \simeq 0.854\,n\sqrt{\lambda n}}$$

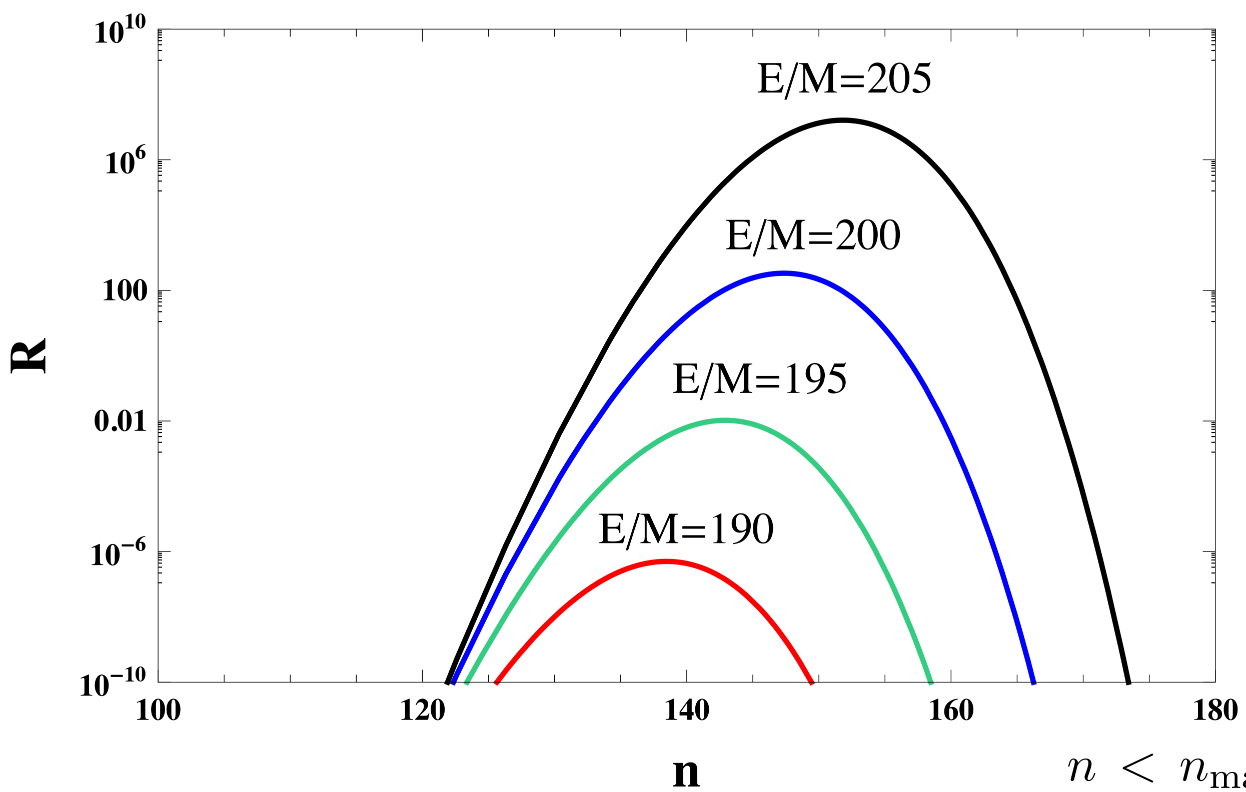


Summary of the semiclassical result

VVK 1705.04365
1806.05648

$$\lambda \to 0\,, \quad n \to \infty\,, \quad \text{with} \quad \lambda n = \text{fixed} \gg 1\,, \quad \varepsilon = \text{fixed} \ll 1$$

$$\mathcal{R}_n(E) = e^{W(E,n)} = \exp\left[\frac{\lambda n}{\lambda}\left(\log\frac{\lambda n}{4} + 0.85\sqrt{\lambda n} - 1 + \frac{3}{2}\left(\log\frac{\varepsilon}{3\pi}+1\right) - \frac{25}{12}\varepsilon\right)\right]$$



$$E/m = (1+\varepsilon)\,n$$

positive
(quantum effects)

negative
(phase space)

Can always make this term win =>
unsuppressed R at high Energies

$n < n_{\max} = E/M_h$

VVK & Spannowsky 1704.0344

Higher order corrections are suppressed by extra powers of $\lambda \to 0$ and $1/n \to 0$ and by $\mathcal{O}(1/\sqrt{\lambda n})$ as well as by $\mathcal{O}(\varepsilon)$.

Thus we have computed the rate R in the large lambda n limit:

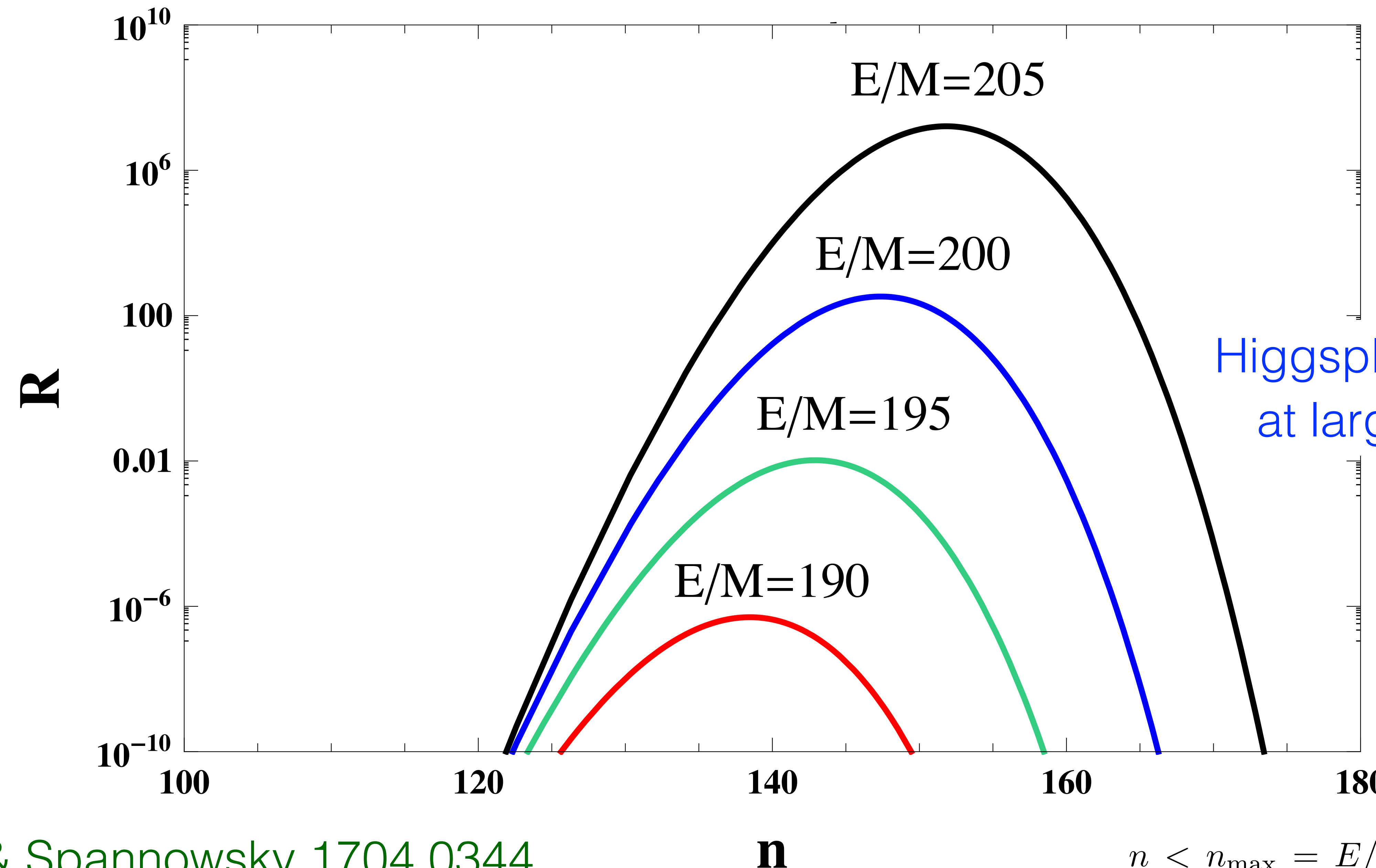
using the semi-classical approach and the thin-wall approximation

VVK 1705.04365

VVK 1806.05648

$$\mathcal{R} = \exp\left[\frac{\lambda n}{\lambda}\left(\log\frac{\lambda n}{4} + \boxed{3.02\sqrt{\frac{\lambda n}{4\pi}}} - 1 + \frac{3}{2}\left(\log\frac{\varepsilon}{3\pi}+1\right) - \frac{25}{12}\varepsilon\right)\right]$$

$\lambda n \gg 1 \quad$ small ε



Higgsplosion realised at large lambda n

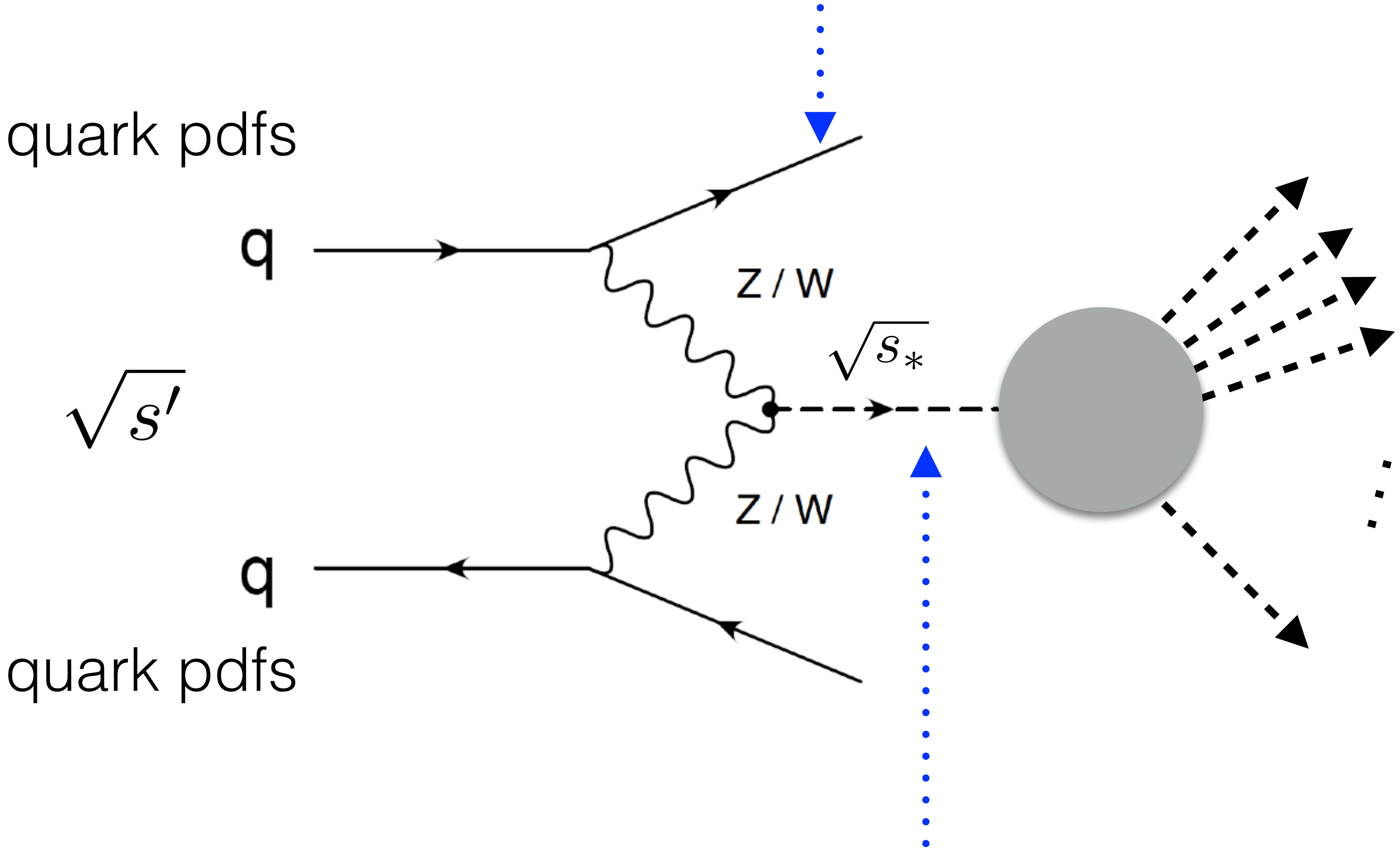
VVK & Spannowsky 1704.0344

$n < n_{\max} = E/M_h$

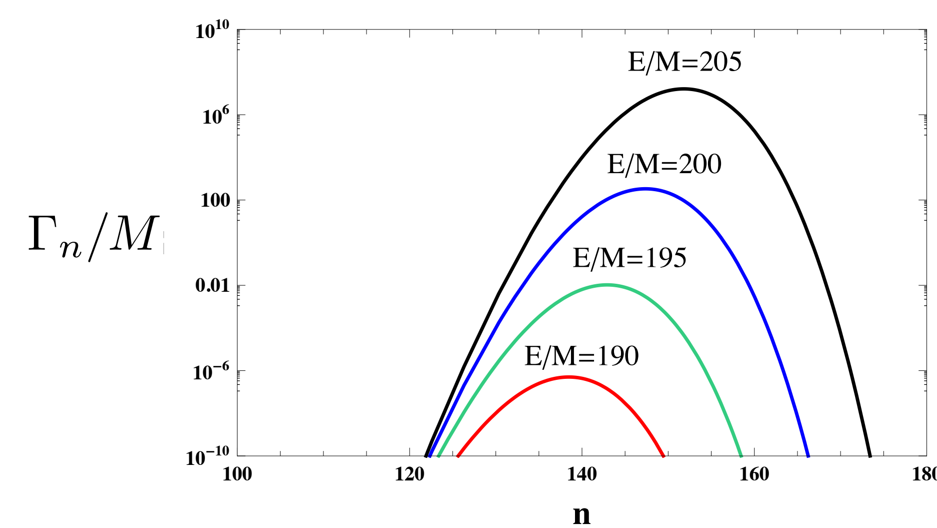
Applications:

Vector boson fusion at high-energy pp colliders (FCC)

energy excess over $\sqrt{s_*}$ carried away by jets



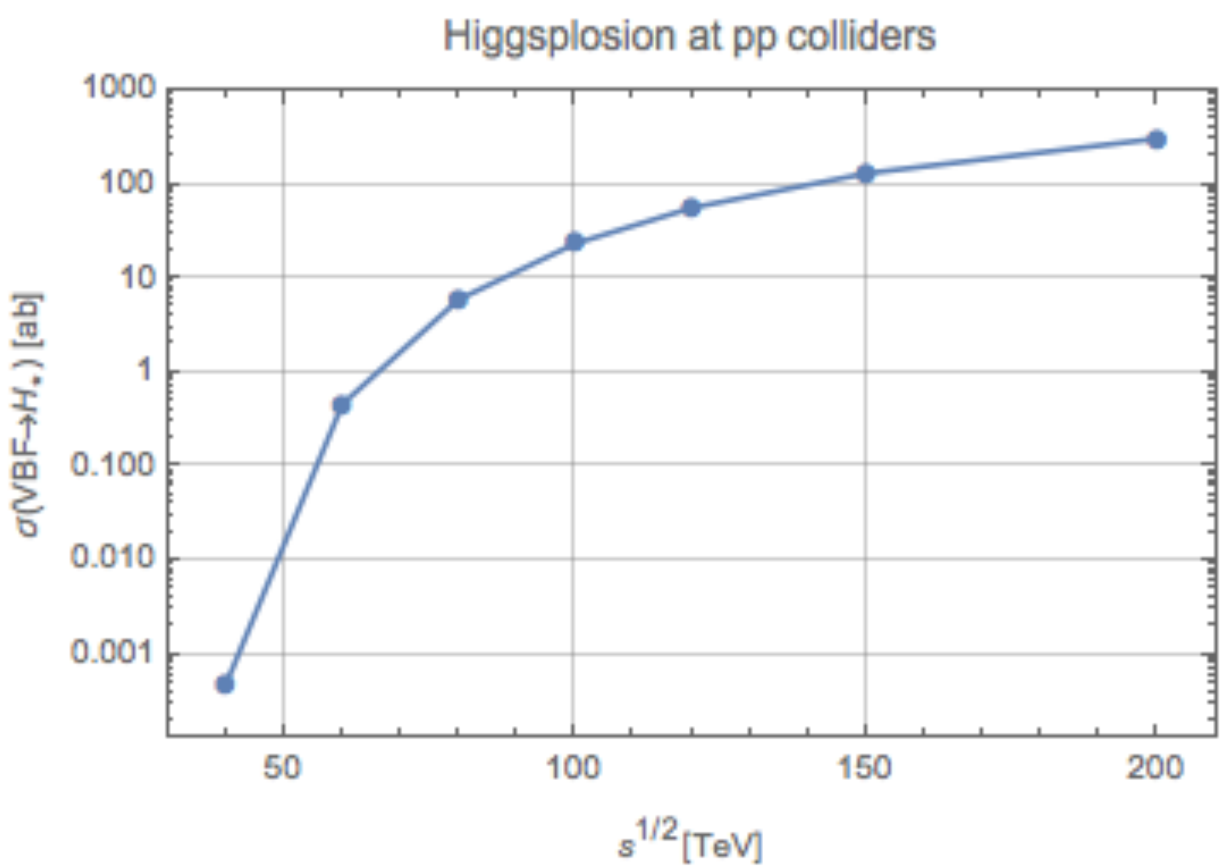
n non-relativistic Higgses
Higgsplosion at $\sqrt{s_*}$

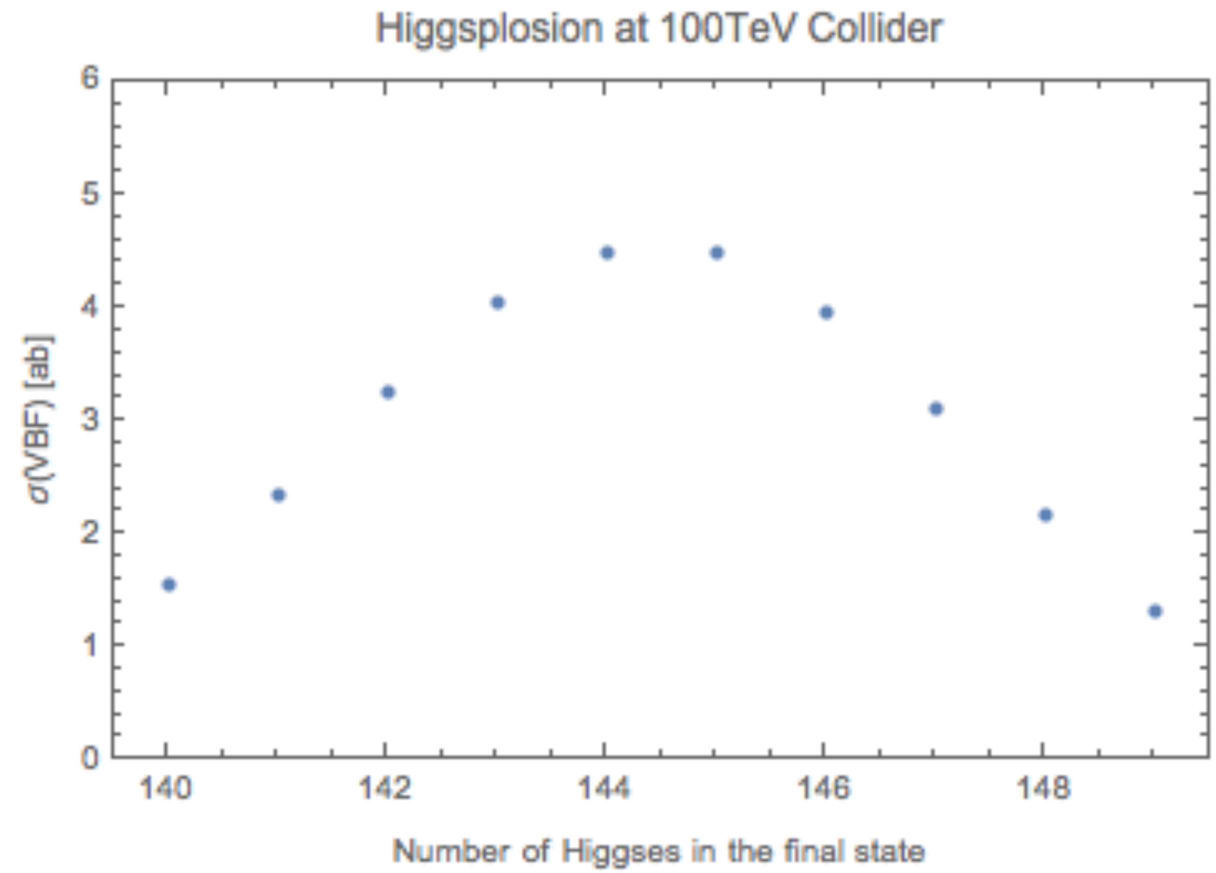


$$\frac{i}{s_* - m_h^2 - Re\tilde{\Sigma}(s_*) + i m_h \Gamma(s_*)}$$

Propagator with Higgspersion at $\sqrt{s_*}$

Vector boson fusion at high-energy pp colliders (FCC)

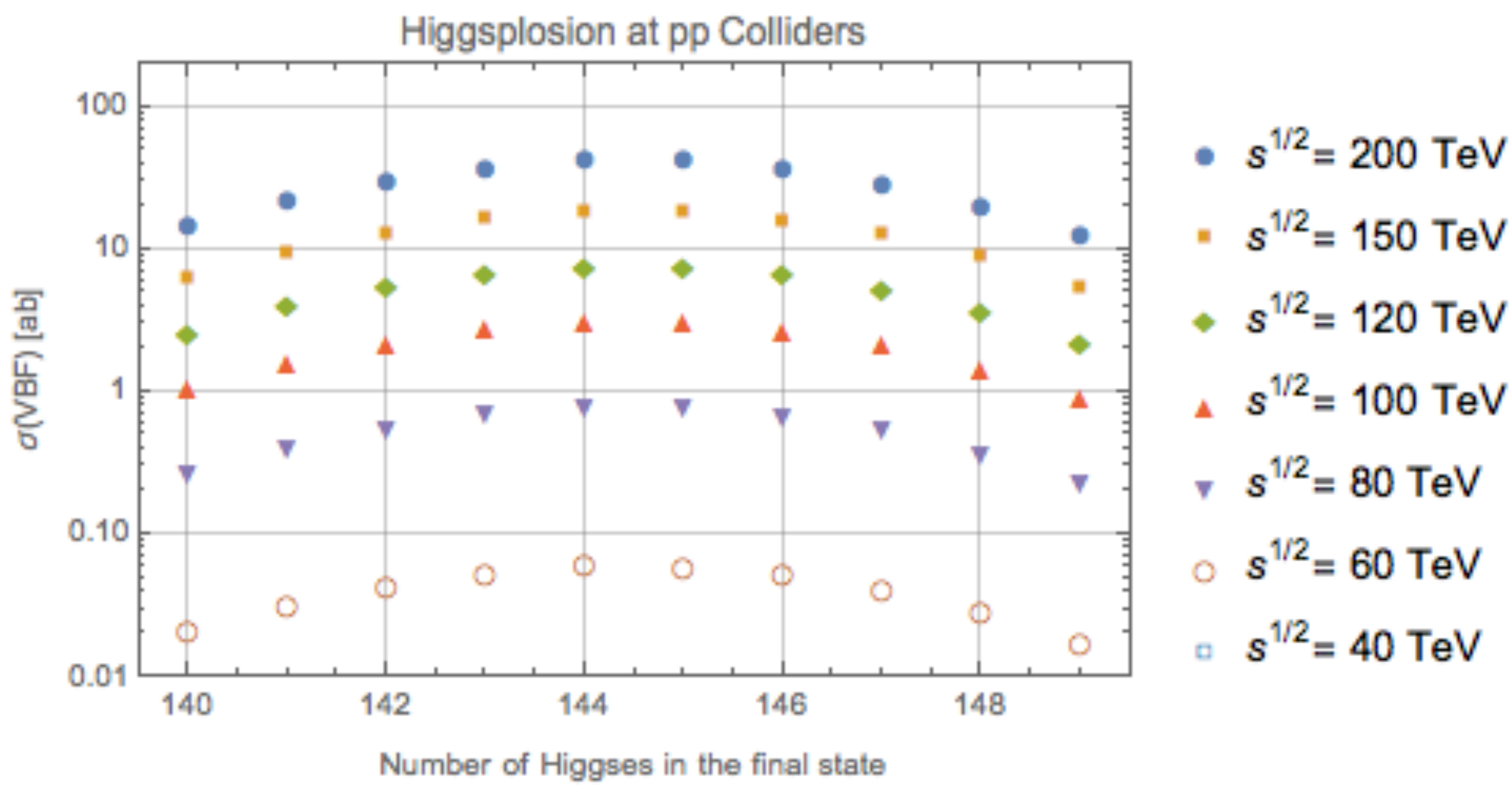




using $p_{t\,\mathrm{jet}} > 40$ GeV

VVK, J Scholtz, M Spannowsky

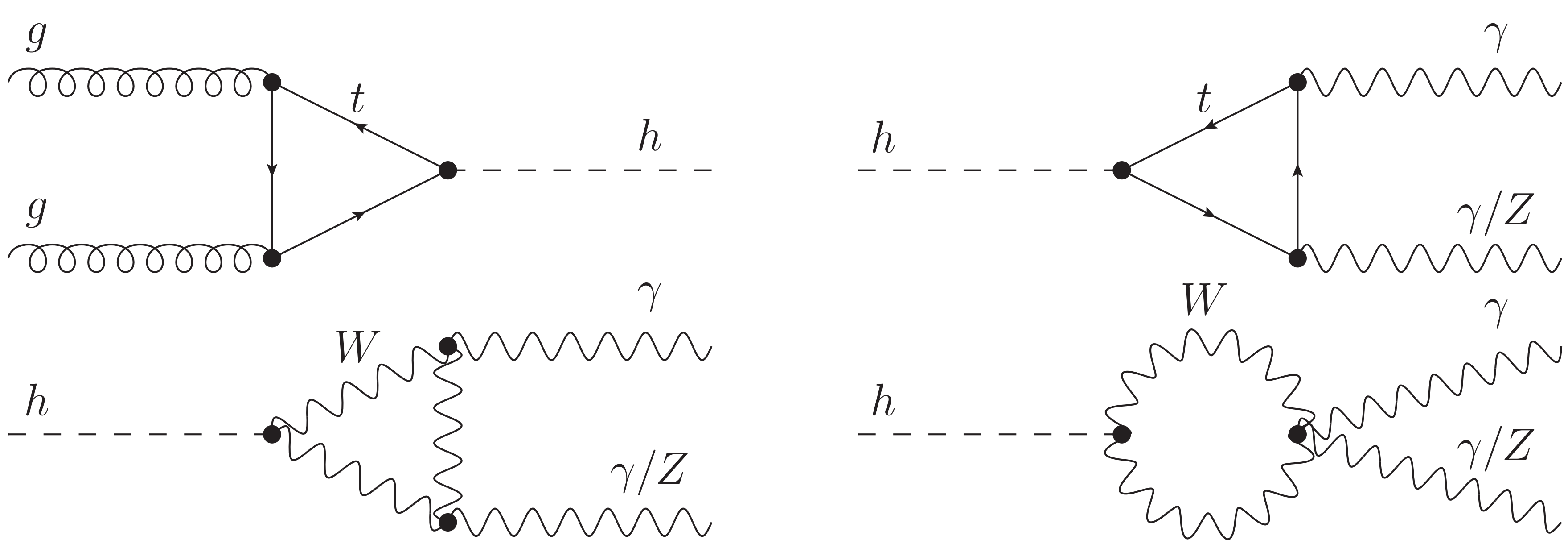
preliminary: no Higgs decays into SM d.o.f included;
& no vector bosons in final states



Effects of Higgsplosion on Precision Observables

- VVK, J Reiness, M Spannowsky, P Waite 1709.08655

Here focus on a class of observables which have no tree-level contributions

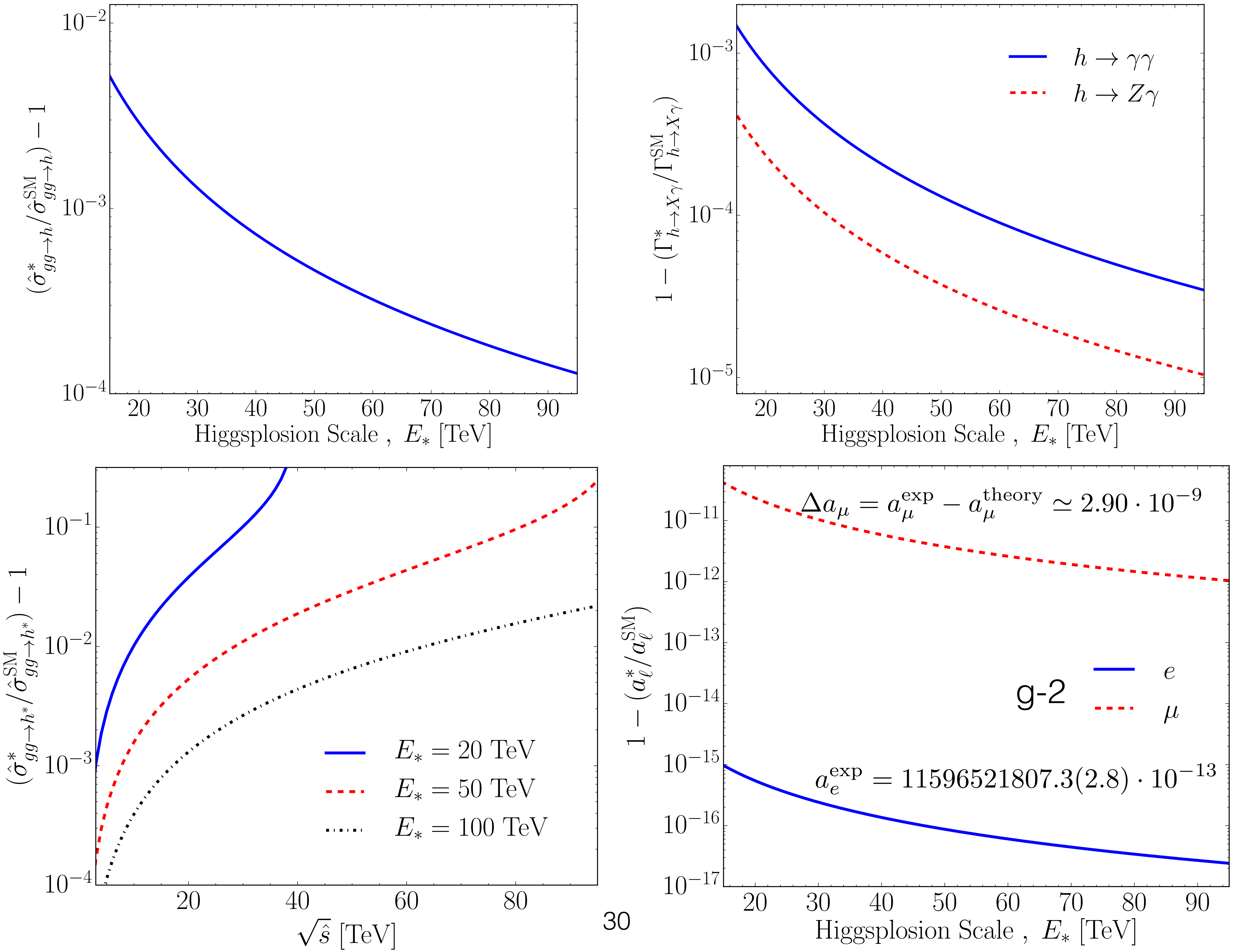


At LHC energies effects of Higgsplosion are small (next slide).

However O(1) effects can be achieved for these loop-induced processes if the interactions are probed close to ~ 2E*.

Effects of Higgsplosion on Precision Observables

VVK, J Reiness, M Spannowsky, P Waite 1709.08655



•

Conclusions:

$$\Delta_R(p) \;=\; \frac{i}{p^2 - m^2 \;-\; \mathrm{Re}\,\Sigma_R(p^2) \;+\; im\Gamma(p^2) \;+ i\epsilon}$$

R $\longleftrightarrow$ Higgsplosion

Loop integrals are effectively cut off at E_* by the exploding width $\Gamma(p^2)$ of the propagating state into the high-multiplicity final states.

The incoming highly energetic state decays rapidly into the multi-particle state made out of soft quanta with momenta $k_i^2 \sim m^2 \lll E_*^2$.

The width of the propagating degree of freedom becomes much greater than its mass: it is no longer a simple particle state.

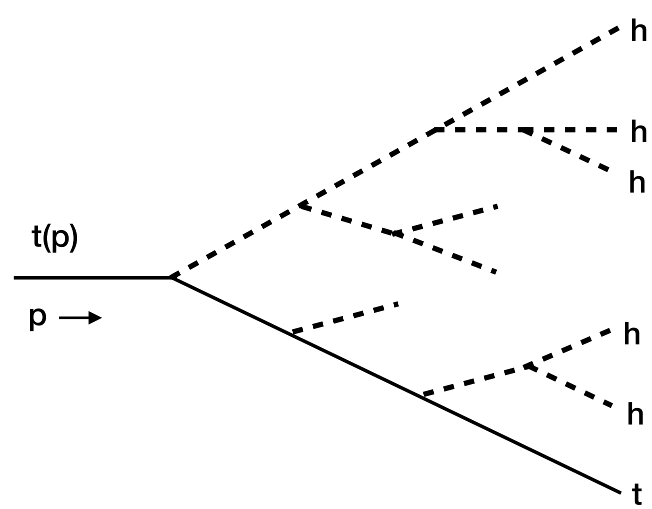
In this sense, it has become a composite state made out of the n soft particle quanta of the same field ϕ.

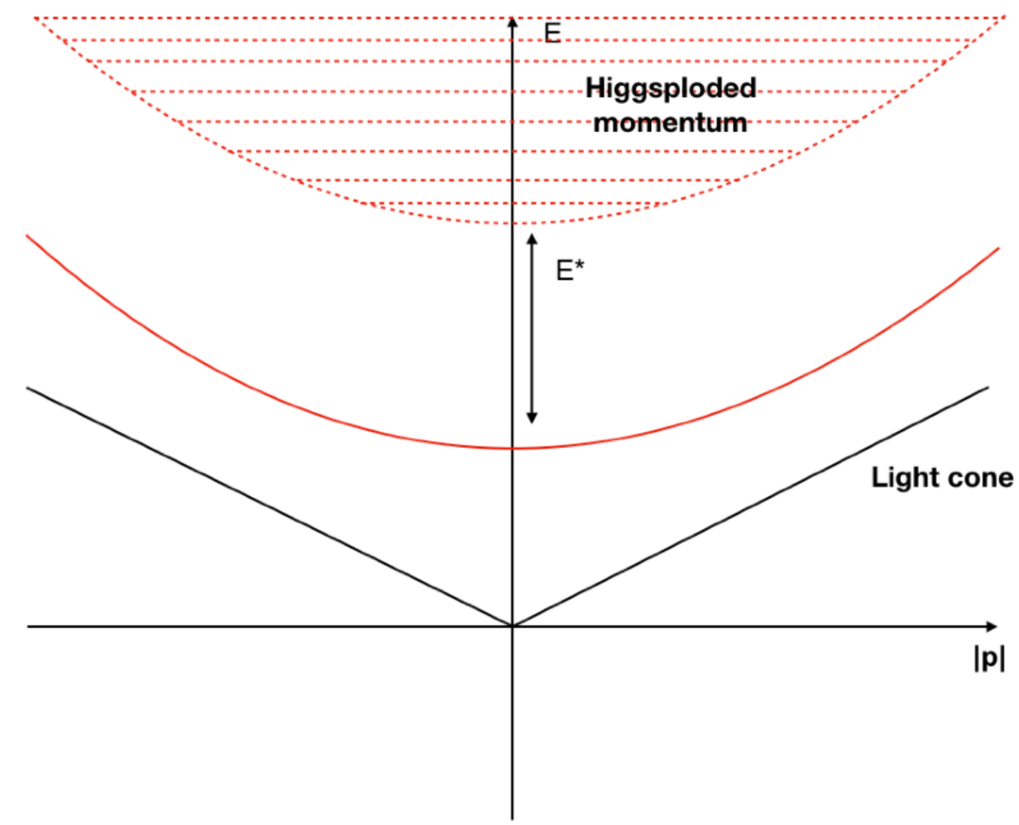
- VVK & Spannowsky 1704.03447, 1707.01531

Consequences of Higgsplosion

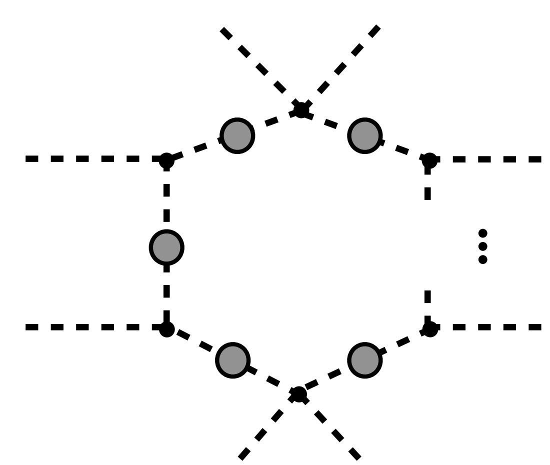
- **All** particles Higgsplode if **virtual** enough

  e.q. top quark, Z boson and even graviton Higgsplodes





- As all virtual particles Higgsplode, virtual corrections are regulated by higgspersing propagators



Consequences of Higgsplosion

- As all loop-diagrams are regulated, i.e. quantum fluctuations are exponentially suppressed, the Standard Model develops an asymptotic fixed point.

  → Classical/Deterministic theory

  → Any highly virtual or a very heavy particle rapidly decays into a large number of relatively soft Higgs bosons. A composite state.

  → Above higgsplosion scale, quantum fluctuations are damped

- SM is embedded into asymptotically safe theory

  → coupling constants stop running above the higgsplosion scale

Consequences of Higgsplosion

- SM has new physical scale

(close analogy to Sphaleron)

$$E_* = C\,\frac{m_h}{\lambda} \quad \text{with} \quad C = \text{const.} \qquad\qquad M_{\rm sph} = \text{const}\,\frac{m_W}{\alpha_w}$$

Scaling behaviour of propagator:

$$\Delta(x) := \langle 0|T(\phi(x)\,\phi(0))|0\rangle \sim \begin{cases} m^2\, e^{-m|x|} & : \text{ for } |x| \gg 1/m \\ 1/|x|^2 & : \text{ for } 1/E_* \ll |x| \ll 1/m \\ E_*^2 & : \text{ for } |x| \lesssim 1/E_* \end{cases},$$

for $|x| \lesssim 1/E_*$ one enters the Higgsplosion regime

$$\frac{i}{s_* - m_h^2 - Re\tilde{\Sigma}(s_*) + i m_h \Gamma(s_*)}$$

Propagator with Higgspersion

Extra slides

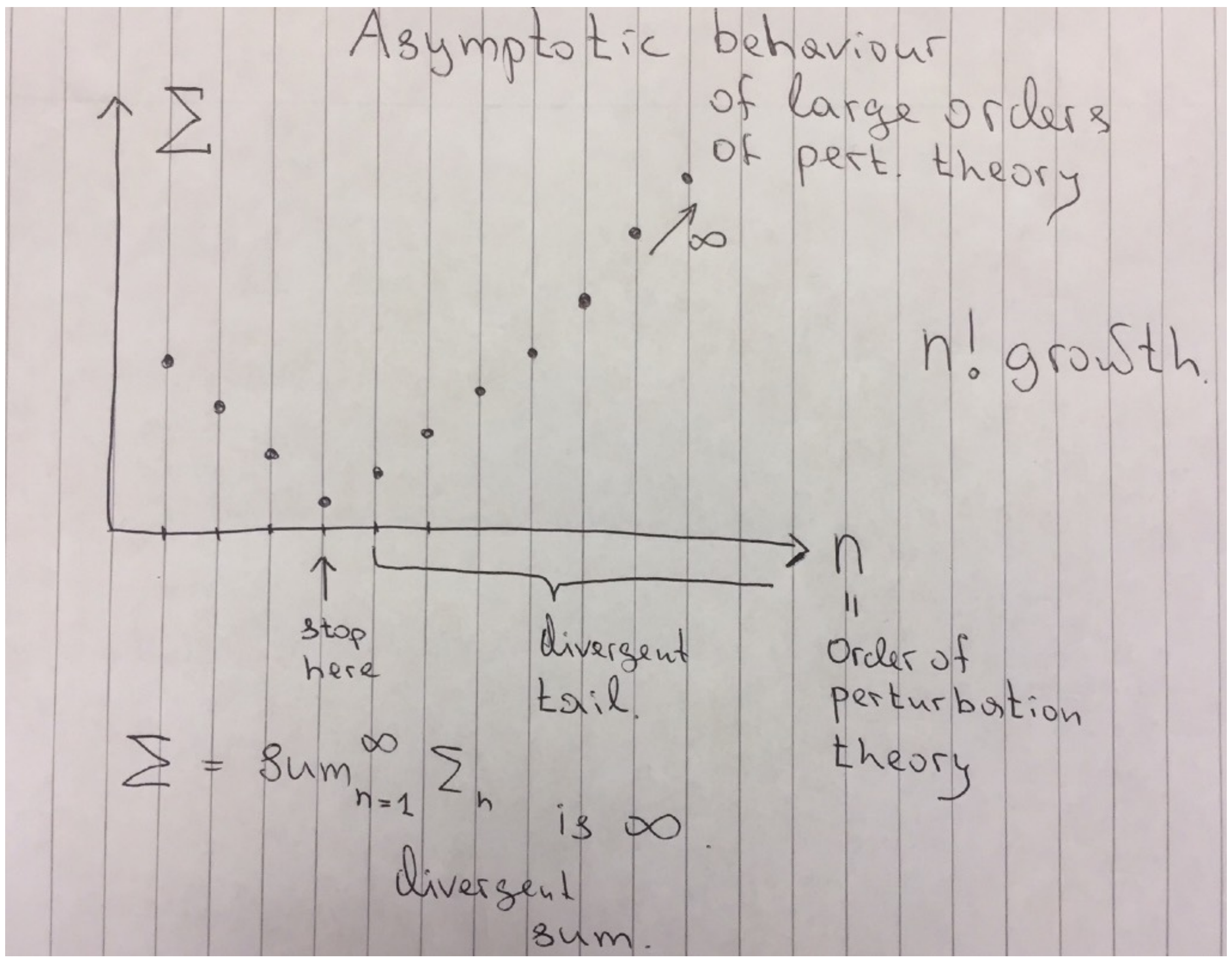
- The n! growth of perturbative amplitudes is not entirely surprising: the number of contributing Feynman diagrams is known to grow factorially with n. [In scalar QFT there are no partial cancellations between individual diagrams (unlike QCD).]

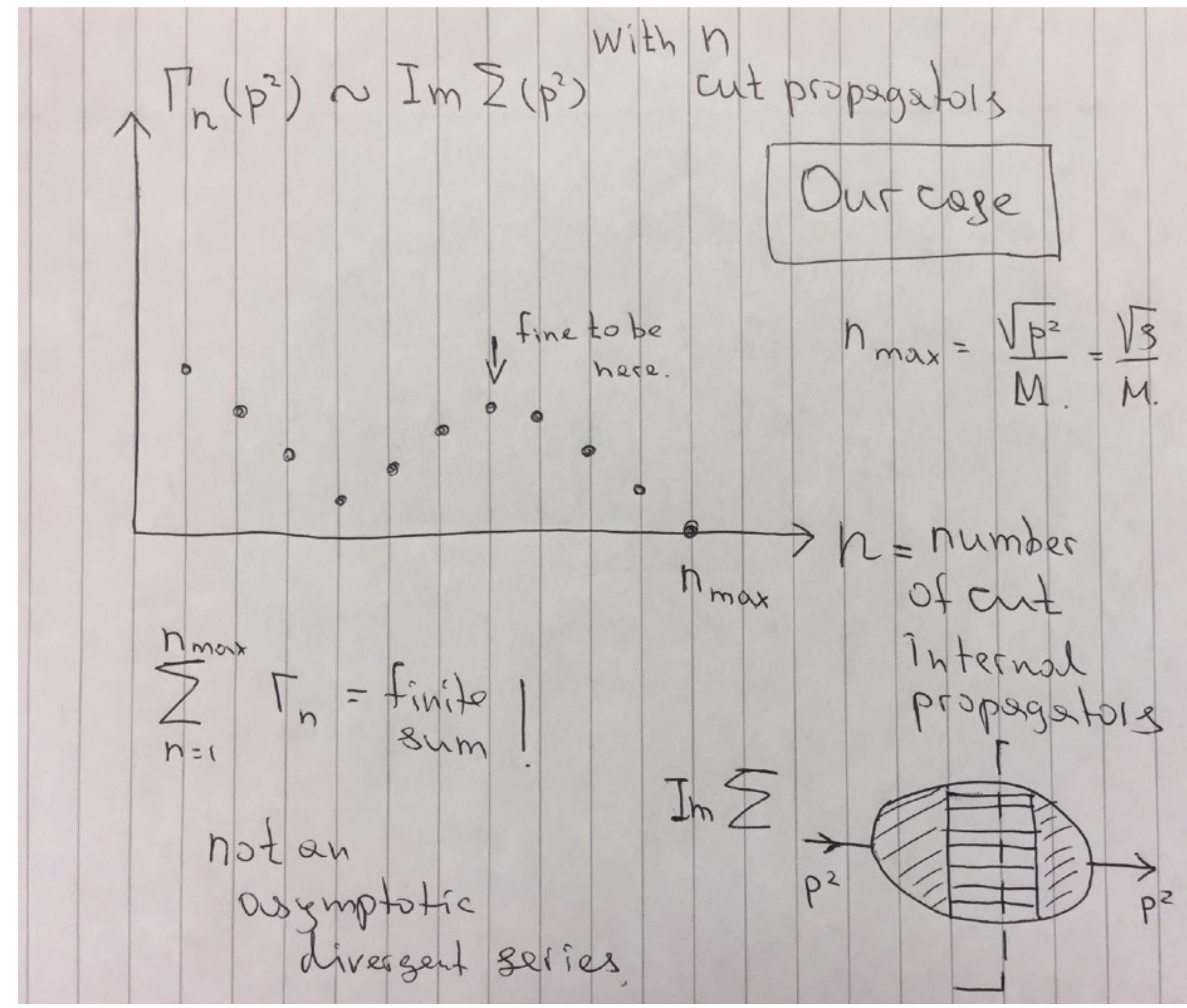
- Important to distinguish between the two types of large-n corrections:

(a) present case where the *leading-order* tree-level contribution to the 1*->n Amplitude grows factorially with the particle multiplicity n of the final state.

- (b) *higher-order* perturbative corrections to some leading-order quantities

Contrast asymptotic growth of higher-order corrections in perturbation theory with the ~n! contributions to Gamma_n(s)





Not the same types of beasts

It is the decay width Gamma_n(s) which is the central object of interest and the driving force of Higgsplosion.