On interference, non-interference, and gauge fixing in the SMEFT

Andreas Helset
Niels Bohr International Academy & Discovery Centre

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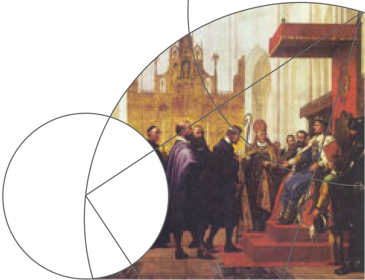

Outline

Based on JHEP **04** (2018) 038 with M. Trott
and arXiv:1803.08001 with M. Paraskevas and M. Trott.

1. Standard Model Effective Field Theory (SMEFT)
2. Gauge fixing
3. Interference and non-interference

Standard Model Effective Field Theory

Using the fields and symmetries of the Standard Model (SM), we add higher dimensional operators

$$\mathcal{L}_{\mathsf{SMEFT}} = \mathcal{L}_{SM} + \mathcal{L}^{(5)} + \mathcal{L}^{(6)} + \mathcal{L}^{(7)} + \dots \qquad (1)$$

where

$$\mathcal{L}^{(d)} = \sum_{i=1}^{n_d} \frac{C_i^{(d)}}{\Lambda^{d-4}} Q_i^{(d)} \qquad \text{for } d > 4. \qquad (2)$$

$C_i^{(d)}$: Wilson coefficient
$Q_i^{(d)}$: Operator with mass dimension d
Λ: Scale of New Physics

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However, using the normal Standard Model gauge fixing procedure leads to Hartmann, Trott, JHEP **07**, 151 (2015)

$$\frac{c_W s_W}{\xi_B \xi_W}(\xi_B - \xi_W)(\partial^\mu A_\mu \partial^\nu Z_\nu) - \frac{C_{HWB} v^2 (s_w^2 - c_W^2)(s_W^2 \xi_B + c_W^2 \xi_W)}{\xi_B \xi_W}(\partial^\mu A_\mu \partial^\nu Z_\nu) + \dots \quad (3)$$

The A-Z mixing doesn't cancel for $\xi_W = \xi_B$.

Geometry of scalar field space

The bilinear field interactions can be thought of in terms of connections on the field space manifold

Alonso, Jenkins, Manohar Phys. Lett. **B754**, 335 (2016), JHEP **08**, 101 (2016)

$$\mathcal{L}_{\text{scalar,kin}} = (D_\mu H)^\dagger (D^\mu H) + \frac{C_{H\Box}}{\Lambda^2} \left(H^\dagger H\right) \Box \left(H^\dagger H\right) + \frac{C_{HD}}{\Lambda^2} \left(H^\dagger D_\mu H\right)^* \left(H^\dagger D^\mu H\right) \tag{4}$$

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where

$$H = \frac{1}{\sqrt{2}} \begin{bmatrix} \phi_2 + i\phi_1 \\ \phi_4 - i\phi_3 \end{bmatrix}, \qquad I, J \in \{1, \dots, 4\}. \tag{5}$$

Metric of the scalar field manifold

The metric is non-trivial

$$h_{IJ}(\phi) = \delta_{IJ} - 2\frac{C_{H\Box}}{\Lambda^2}\phi_I\phi_J + \frac{1}{2}\frac{C_{HD}}{\Lambda^2}f_{IJ}(\phi), \qquad (6)$$

where

$$f_{IJ}(\phi) = \begin{bmatrix} a & 0 & d & c \\ 0 & a & c & -d \\ d & c & b & 0 \\ c & -d & 0 & b \end{bmatrix}, \quad \begin{aligned} a &= \phi_1^2 + \phi_2^2 \\ b &= \phi_3^2 + \phi_4^2 \\ c &= \phi_1\phi_4 + \phi_2\phi_3, \\ d &= \phi_1\phi_3 - \phi_2\phi_4. \end{aligned} \qquad (7)$$

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- The Riemann curvature tensor calculated from the scalar field metric is non-vanishing. The scalar manifold is curved due to the power counting expansion.

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- The Riemann curvature tensor calculated from the scalar field metric is non-vanishing. The scalar manifold is curved due to the power counting expansion.
- Field redefinitions cannot turn the metric into a trivial form.
- Physical quantities depend on field redefinition invariant quantities.

Geometry of gauge field space

Analogously, we can describe the kinetic part of the gauge fields in terms of connections on the field space manifold

$$\mathcal{L}_{\text{gauge,kin}} = -\frac{1}{4}W^a_{\mu\nu}W^{a,\mu\nu} - \frac{1}{4}B_{\mu\nu}B^{\mu\nu} + \frac{C_{HB}}{\Lambda^2}H^\dagger H B_{\mu\nu}B^{\mu\nu} + \frac{C_{HW}}{\Lambda^2}H^\dagger H W^a_{\mu\nu}W^{a,\mu\nu} + \frac{C_{HWB}}{\Lambda^2}H^\dagger \sigma^a H W^a_{\mu\nu}B^{\mu\nu} \quad (8)$$

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where

$$\begin{aligned} g_{ab} &= \left(1 - 4\frac{C_{HW}}{\Lambda^2}H^\dagger H\right)\delta_{ab}, \qquad g_{44} = 1 - 4\frac{C_{HB}}{\Lambda^2}H^\dagger H, \\ g_{a4} &= g_{4a} = -2\frac{C_{HWB}}{\Lambda^2}H^\dagger \sigma_a H, \qquad a = 1, 2, 3. \end{aligned} \tag{9}$$

The gauge field manifold is curved.

Background field method

- The background field method splits fields into background and quantum fields $F \to \hat{F} + F$.
  $\hat{F}$: background field
  F: quantum field

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  $\hat{F}$: background field
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- The background field method provides technical simplifications due to the background field gauge invariance being preserved and the resulting Ward identities.
- The Standard Model was formulated using the background field method Denner, Dittmaier, Weiglein Nucl. Phys. **B440**, 95 (1995)

Gauge fixing the Standard Model

Using the background field method, the electroweak Standard Model gauge fixing term takes the form

Denner, Dittmaier, Weiglein Nucl. Phys. **B440**, 95 (1995)

$$\begin{aligned}\mathcal{L}_{\mathrm{GF}} = &-\frac{1}{2\xi_W}\sum_{a=1}^{3}\Bigg[\partial_\mu W^{a,\mu} - g_2\epsilon^{abc}\hat{W}_{b,\mu}W_c^\mu \\ &+ig_2\frac{\xi_W}{2}\left(\hat{H}_i^\dagger(\sigma^a)^i_j H^j - H\dagger_i(\sigma^a)^i_j\hat{H}^j\right)\Bigg]^2 \\ &-\frac{1}{2\xi_B}\left[\partial_\mu B^\mu + ig_1\frac{\xi_B}{2}\left(\hat{H}_i^\dagger H^i - H_i^\dagger\hat{H}^i\right)\right]^2 .\end{aligned} \tag{10}$$

$\hat{W}, \hat{B}, \hat{H}$: background fields
W, B, H: quantum fields.
Background field gauge invariance is preserved.

Real representation of the scalar field

We cannot use the Pauli matrix representation when we have the ϕ^I fields. We use the real representation

$$\gamma^I_{1,J} = \begin{bmatrix} 0 & 0 & 0 & -1 \\ 0 & 0 & -1 & 0 \\ 0 & 1 & 0 & 0 \\ 1 & 0 & 0 & 0 \end{bmatrix} \quad \gamma^I_{2,J} = \begin{bmatrix} 0 & 0 & 1 & 0 \\ 0 & 0 & 0 & -1 \\ -1 & 0 & 0 & 0 \\ 0 & 1 & 0 & 0 \end{bmatrix}$$

$$\gamma^I_{3,J} = \begin{bmatrix} 0 & -1 & 0 & 0 \\ 1 & 0 & 0 & 0 \\ 0 & 0 & 0 & -1 \\ 0 & 0 & 1 & 0 \end{bmatrix} \quad \gamma^I_{4,J} = \begin{bmatrix} 0 & -1 & 0 & 0 \\ 1 & 0 & 0 & 0 \\ 0 & 0 & 0 & 1 \\ 0 & 0 & -1 & 0 \end{bmatrix}. \tag{11}$$

We have that

$$[\gamma_a, \gamma_b] = 2\epsilon^c_{ab}\gamma_c, \qquad \tilde{\gamma}_A = \begin{cases} g_2\gamma_A & \text{for } A = 1,2,3 \\ g_1\gamma_A & \text{for } A = 4, \end{cases}$$

$$[\gamma_a, \gamma_4] = 0, \qquad \tilde{\epsilon}^A_{BC} = g_2\epsilon^A_{BC}. \tag{12}$$

Gauge fixing the Standard Model Effective Field Theory

A gauge fixing choice which preserves the geometric structure of the theory is

$$\mathcal{L}_{\rm GF} = -\frac{\hat{g}_{AB}}{2\xi}\mathcal{G}^A\mathcal{G}^B,$$
$$\mathcal{G}^X \equiv \partial_\mu W^{X,\mu} - \tilde{\epsilon}^X_{CD}\hat{W}^C_\mu W^{D,\mu} + \frac{\xi}{2}\hat{g}^{XC}\phi^I\hat{h}_{IK}\tilde{\gamma}^K_{C,J}\hat{\phi}^J. \quad (13)$$

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Background field gauge transformations

It is useful to note the following background field gauge transformations ($\delta\hat{F}$), with infinitesimal local gauge parameters $\delta\hat{\alpha}_A(x)$ when verifying the explicitly the background field gauge invariance of this expression

$$
\begin{aligned}
\delta\,\hat{\phi}^I &= -\delta\hat{\alpha}^A\,\frac{\tilde{\gamma}^I_{A,J}}{2}\hat{\phi}^J,\\
\delta\,(D^\mu\hat{\phi})^I &= -\delta\hat{\alpha}^A\,\frac{\tilde{\gamma}^I_{A,J}}{2}(D^\mu\hat{\phi})^J,\\
\delta\hat{\mathcal{W}}^{A,\mu} &= -\partial^\mu(\delta\hat{\alpha}^A) - \tilde{\epsilon}^A{}_{BC}\,\delta\hat{\alpha}^B\,\hat{\mathcal{W}}^{C,\mu},\\
\delta\hat{h}_{IJ} &= \hat{h}_{KJ}\,\frac{\delta\hat{\alpha}^A\,\tilde{\gamma}^K_{A,I}}{2} + \hat{h}_{IK}\,\frac{\delta\hat{\alpha}^A\,\tilde{\gamma}^K_{A,J}}{2},\\
\delta\hat{\mathcal{W}}^A_{\mu\nu} &= -\tilde{\epsilon}^A{}_{BC}\,\delta\hat{\alpha}^B\,\hat{\mathcal{W}}^C_{\mu\nu},\\
\delta\hat{g}_{AB} &= \hat{g}_{CB}\,\tilde{\epsilon}^C{}_{DA}\,\delta\hat{\alpha}^D + \hat{g}_{AC}\,\tilde{\epsilon}^C{}_{DB}\,\delta\hat{\alpha}^D.
\end{aligned}
\tag{14}
$$

The background field gauge invariance is established by using these transformations in conjuction with a linear change of variables on the quantum fields

$$\mathcal{W}^{A,\mu} \to \mathcal{W}^{A,\mu} - \tilde{\epsilon}^{A}_{\ \,BC}\, \delta\hat{\alpha}^{B}\, \mathcal{W}^{C,\mu},$$
$$\phi^{I} \to \phi^{I} - \frac{\delta\hat{\alpha}^{B}\, \tilde{\gamma}^{I}_{B,K}}{2} \phi^{K}. \tag{15}$$

The transformation of the gauge fixing term is

$$\delta\mathcal{G}^{X} = -\tilde{\epsilon}^{X}_{\ \,AB}\, \delta\hat{\alpha}^{A}\mathcal{G}^{B}. \tag{16}$$

With these transformations, the background field gauge invariance of the gauge fixing term is directly established.

Ghost term

The quantum fields gauge transformations are

$$\Delta\mathcal{W}^A_\mu = -\partial_\mu \Delta\alpha^A - \tilde{\epsilon}^A{}_{BC}\,\Delta\alpha^B\,(\mathcal{W}^C_\mu + \hat{\mathcal{W}}^C_\mu),$$
$$\Delta\phi^I = -\Delta\alpha^A\,\frac{\tilde{\gamma}^I_{A,J}}{2}\,(\phi^J + \hat{\phi}^J). \tag{17}$$

As the hatted field metrics depend only on the background fields and do not transform under quantum field gauge transformations, the Faddeev-Popov ghost term still follows directly; we find

$$\begin{aligned}\mathcal{L}_{\rm FP} = &- \hat{g}_{AB}\bar{u}^B\Big[-\partial^2\delta^A_C - \overleftarrow{\partial}_\mu\tilde{\epsilon}^A{}_{DC}(\mathcal{W}^{D,\mu} + \hat{\mathcal{W}}^{D,\mu})\\ &+ \tilde{\epsilon}^A{}_{DC}\hat{\mathcal{W}}^D_\mu\overrightarrow{\partial}^\mu - \tilde{\epsilon}^A{}_{DE}\tilde{\epsilon}^E{}_{FC}\hat{\mathcal{W}}^D_\mu(\mathcal{W}^{F,\mu} + \hat{\mathcal{W}}^{F,\mu})\\ &- \frac{\xi}{4}\hat{g}^{AD}(\phi^J + \hat{\phi}^J)\tilde{\gamma}^I_{C,J}\,\hat{h}_{IK}\,\tilde{\gamma}^K_{D,L}\,\hat{\phi}^L)\Big]\,u^C.\end{aligned} \tag{18}$$

Summary

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The key point

We gauge fix the fields on the curved field space due to the power counting expansion.

Interference in the SMEFT

When the Standard Model (SM) and Beyond Standard Model (BSM) contribute to the same amplitude

$$A_{\mathrm{SMEFT}} = A_{\mathrm{SM}} + \frac{1}{\Lambda^2} A^{(6)} + \dots \tag{19}$$

where $A^{(6)}$ is an amplitude with one insertion of an operator from $\mathcal{L}^{(6)}$. The cross section

$$\sigma \propto |A_{\mathrm{SMEFT}}|^2 \sim |A_{\mathrm{SM}}|^2 + \frac{1}{\Lambda^2} A_{\mathrm{SM}} \times A^{(6)} + \frac{1}{\Lambda^4} |A^{(6)}|^2 + \dots \tag{20}$$

For small BSM effects the interference term dominates over the last term.

Non-interference in the SMEFT

When the SM and BSM do not contribute to the same amplitude

$$\sigma \propto \sum |A_{\mathrm{SMEFT}}|^2 \sim |A_{\mathrm{SM}}|^2 + \frac{1}{\Lambda^4}|A^{(6)}|^2 + \ldots \qquad (21)$$

The leading BSM effects are suppressed by $\mathcal{O}(\frac{1}{\Lambda^4})$, the same order as operators of mass dimension 8 that do interfere.

Hard to measure!

More on non-inteference

The phenomenon of non-interference has been seen before in a QCD context (Simmons '89, Dixon and Shadmi '94).
General statements can be made from helicity arguments:

Non-interference statement

Four-point amplitudes with at least one transverse polarized gauge boson do not interfere at tree level in the massless limit.

Lately, similar reasoning has been applied to electroweak diboson production in the high energy limit (Azatov et. al. '16).

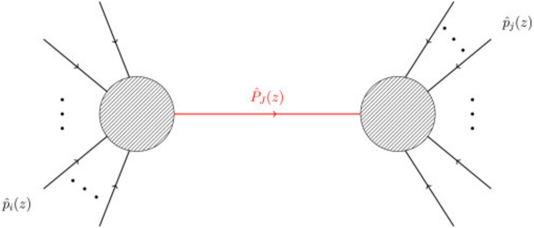
Helicity arguments

For massless particles we define helicity.

We put intermediate propagators on-shell:

$$h(A_n) = h(A_m) + h(A'_m) \tag{22}$$

where $n = m + m' - 2$.



Helicity arguments

Using little group scaling and dimensional analysis, we have that

$$|h(A_3)| = 1 - [g] \tag{23}$$

In the SM, $|h_3^{\text{SM}}| = 1$, while for dimension-6 operators $|h_3^{\text{BSM}}| = 3$. For the SM at we can use helicity selection rules (from SUSY Ward identities) to show that

$$|h(A_4)^{\text{SM}}| < 2 \qquad \text{(at least one vector boson)} \tag{24}$$

Helicity arguments

Summary:

- Helicity sums
- Three-point kinematics
- Helicity selection rules

Result:

$$|h_4^{\text{SM}}| = 0 \tag{25}$$
$$|h_4^{\text{BSM}}| = 2, 4 \tag{26}$$

when there is at least one transverse vector boson.

Softening the claim

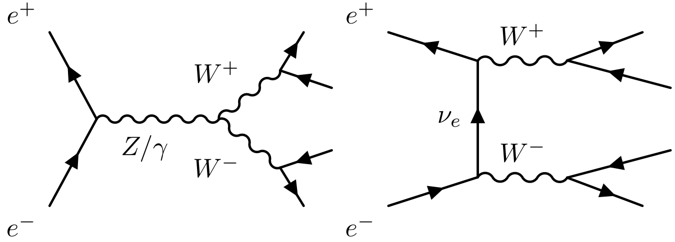
For electroweak diboson production we note that it holds:

- Only at tree level
- In the high energy limit, $\hat{m}_W^2/s \ll 1$
- For on-shell vector bosons

The statement will get loop and mass corrections.
In addition, on-shell massive gauge bosons are not formally physical.

The first two points have been considered.
We also take the last possibility into account.

Off-shell effects



We investigate three regions of phase space:

- Case 1: both $W^\pm$ near on-shell
- Case 2: both $W^\pm$ off-shell
- Case 3: one $W^\pm$ near on-shell

Only Case 1 has non-interference. Off-shell effects are suppressed by the width.

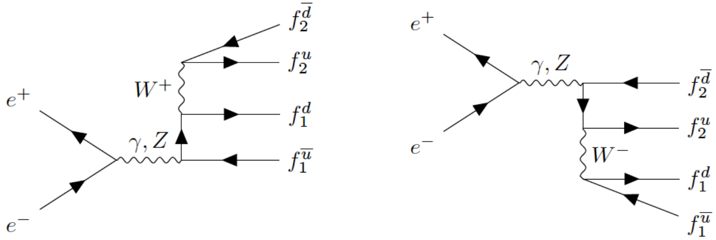
Off-shell effects

$\lambda_{12}\lambda_{34}\lambda_+\lambda_-$	$\sum_X M_X^{\lambda^\pm}/4\pi\hat{\alpha}$
$00-+$	$\frac{\sin\theta}{2\sqrt{s_1 s_3}}\left[\frac{1}{c_{\hat{\theta}}^2}+\left(\delta\kappa^{Z\alpha}-\delta F_2^{Z\alpha}\right)y\right]$
$\pm\pm-+$	$-\sin\theta\left[\frac{x^2}{c_{\hat{\theta}}^2}+\frac{y\,\delta\lambda^{Z\alpha}}{2}+\left(\delta g_1^{Z\alpha}-\delta F_2^{Z\alpha}-(s_1+s_3)\frac{\lambda_{Z\alpha}}{2}+\frac{\delta\lambda_Z}{2\,c_{\hat{\theta}}^2}\right)y\,x^2\right]$
$\pm0-+$	$-\frac{(1\pm\cos\theta)\,x}{\sqrt{2s_3}}\left[\frac{1}{c_{\hat{\theta}}^2}+\frac{y}{2}\left(\delta g_1^{Z\alpha}-2\delta F_2^{Z\alpha}+\delta\kappa^{Z\alpha}+s_3\,\delta\lambda^{Z\alpha}\right)\right]$
$0\pm-+$	$\frac{(1\mp\cos\theta)\,x}{\sqrt{2s_1}}\left[\frac{1}{c_{\hat{\theta}}^2}+\frac{y}{2}\left(\delta g_1^{Z\alpha}-2\delta F_2^{Z\alpha}+\delta\kappa^{Z\alpha}+s_1\delta\lambda^{Z\alpha}\right)\right]$
$00+-$	$\frac{\sin\theta}{2\sqrt{s_1 s_3}}\left[(\frac{s_{\hat{\theta}}^2-c_{\hat{\theta}}^2}{2\,c_{\hat{\theta}}^2\,s_{\hat{\theta}}^2})+\frac{s_1+s_3}{2\,s_{\hat{\theta}}^2}+\left(\delta\kappa^{Z\alpha}-\frac{\delta\kappa_Z}{2\,s_{\hat{\theta}}^2}+\frac{2\,\delta\bar{g}_W^\ell}{s_{\hat{\theta}}^2}-\delta F_2^{Z\alpha}\right)y\right]$
$\pm\pm+-$	$-\frac{\sin\theta}{2}\left[\left(1-\frac{1}{2s_{\hat{\theta}}^2}\right)\delta\lambda_Z-\delta\lambda_\alpha\right]y$
$\pm0+-$	$\frac{(1\mp\cos\theta)\,x}{2\sqrt{2s_3}}\left[(\frac{s_{\hat{\theta}}^2-c_{\hat{\theta}}^2}{c_{\hat{\theta}}^2\,s_{\hat{\theta}}^2})+\frac{s_1}{s_{\hat{\theta}}^2}+\frac{s_3}{s_{\hat{\theta}}^2}\frac{1\pm2+3\cos\theta}{1+\cos\theta}-y\frac{\left(\delta g_1^z+\delta\kappa_z+s_3\,\delta\lambda_z\right)}{2s_{\hat{\theta}}^2}\right.$
	$\left.-y\left(\delta F_1^{Z\alpha}-\frac{4\,\delta\bar{g}_W^\ell}{s_{\hat{\theta}}^2}-\left(\delta g_1^{Z\alpha}+\delta\kappa^{Z\alpha}+s_3\,\delta\lambda^{Z\alpha}\right)\right)\right]$
$0\pm+-$	$-\frac{(1\pm\cos\theta)\,x}{2\sqrt{2s_1}}\left[(\frac{s_{\hat{\theta}}^2-c_{\hat{\theta}}^2}{c_{\hat{\theta}}^2\,s_{\hat{\theta}}^2})+\frac{s_3}{s_{\hat{\theta}}^2}+\frac{s_1}{s_{\hat{\theta}}^2}\frac{1\mp2+3\cos\theta}{1+\cos\theta}-y\frac{\left(\delta g_1^z+\delta\kappa_z+s_1\,\delta\lambda_z\right)}{2s_{\hat{\theta}}^2}\right.$
	$\left.-y\left(\delta F_1^{Z\alpha}-\frac{4\,\delta\bar{g}_W^\ell}{s_{\hat{\theta}}^2}-\left(\delta g_1^{Z\alpha}+\delta\kappa^{Z\alpha}+s_1\,\delta\lambda^{Z\alpha}\right)\right)\right]$
$\pm\mp+-$	$\frac{(\mp1+\cos\theta)\sin\theta}{2\,s_{\hat{\theta}}^2(1+\cos\theta)}$

Table 1: Expansion in $x, y < 1$ for the near on-shell region of phase space of the CC03 diagrams approximating $\bar{\psi}\psi \to \bar{\psi}'_1\psi'_2\bar{\psi}'_3\psi'_4$. For exactly on-shell intermediate $W^\pm$ bosons $s_1 = s_3 = 1$. We have used the notation $\delta F^i_{Z\alpha} = (\delta F_i^Z + \delta F_i^\alpha)/4\pi\hat{\alpha}$, $\delta\lambda^{Z\alpha} = \delta\lambda_Z - \delta\lambda_\alpha$, $\delta\kappa^{Z\alpha} = \delta\kappa_Z - \delta\kappa_\alpha$ and $\delta g_1^{Z\alpha} = \delta g_1^Z - \delta g_1^\alpha$.

Gauge invariance

To ensure gauge invariance, we include single resonant diagrams



This does not affect the other results.

Summary

- For electroweak diboson production the SM and the SMEFT interference vanishes in on-shell regions of phase space
- However, for off-shell regions of phase space, interference is restored