SCREENING THEOREM FOR THE HIGGS DECAYS INTO THE
GAUGE BOSON PAIRS

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ABSTRACT

The radiative corrections to the decays of the neutral CP-even Higgs boson $H$ into a longitudinal gauge boson pair, i.e., $H \rightarrow Z_L Z_L$ and $W_L^+ W_L^-$ are analyzed in the two Higgs doublet model by making use of the equivalence theorem. The sensitivity of the decay rates to the masses of the heavier Higgs bosons, charged $G^\pm$ and CP-odd neutral $A$ bosons as well as CP-even neutral $h$ boson, is investigated. Though the width $\Gamma(H \rightarrow Z_L Z_L)$ is insensitive to the masses of heavier Higgs bosons, $\Gamma(H \rightarrow W_L^+ W_L^-)$ is sensitive and the radiative corrections are minimized for $m_G = m_A$. These results are explained completely on the basis of a new screening theorem for the vertices, which is closely connected with the custodial $SU(2)_V$ symmetry.

I Introduction

The discovery of the Higgs boson is the most important task for future accelerators including the next linear colliders. Let us suppose that the Higgs boson $H$ would be discovered in these accelerators and that the mass would turn out to be above the threshold

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of a gauge boson pair. The Higgs sector is then expected to be a strong-coupled sector. The decay of the $H$ boson is then dominated by the decay processes into the longitudinal gauge boson pair, $H \to W_L^+W_L^-$ and $Z_LZ_L$, and these decay widths would be investigated experimentally to a considerable extent. Since studies on the non-decoupling effects due to internal particles in the radiative corrections are expected to be helpful to the approach to the strong Higgs sector, we have calculated the radiative corrections to these decay rates in the two Higgs doublet model (THDM) and have investigated the virtual effects on these due to the internal heavy scalar bosons, charged $G^\pm$, CP-odd neutral $A$ and the another CP-even neutral $h$ boson. These results are summarized in (I) and (II) below.

(I) The radiative corrections to the decay width $\Gamma(H \to W_L^+W_L^-)$ are sensitive to the masses of $G^\pm$ and $A$ bosons ($m_G$ and $m_A$, respectively) but insensitive to the mass of $h$ boson and are minimized if we set $m_G = m_A$.

(II) The radiative corrections to the width $\Gamma(H \to Z_LZ_L)$ are relatively small and are insensitive to the masses of all the internal heavy scalar bosons.

The main purpose of this talk is rather to present a new screening theorem:

**Theorem:** In the radiative corrections to these decay widths the leading contributions with respect to the masses of the internal heavy scalar bosons ($G^\pm$, $A$ and $h$) cancel out in the custodial $SU(2)_V$ symmetric limit of the model.

The characteristic features of the decay widths mentioned in (I) and (II) can be given satisfactory explanations in terms of this theorem. This theorem reminds us of the Veltman’s screening theorem [1]: the leading contributions of the Higgs boson mass in the oblique-type corrections cancel out by virtue of the custodial $SU(2)_V$ symmetry which becomes exact in the weak $U(1)_Y$-coupling limit in the standard model with one Higgs doublet. Veltman’s theorem has been proved by Einhorn and Wudka [2] to all orders. In view of the similarity between Veltman’s theorem for oblique-type corrections and our
counterpart for vertices, our theorem is likely to hold to all orders of the perturbation, though we confirm this theorem only at one loop level.

In our calculation, since we assume that $H$ boson is the lightest of all the Higgs bosons and is much heavier than the gauge bosons, we will make good use of the equivalence theorem. The calculations in the one Higgs-doublet model from the same viewpoint as ours are seen in ref.[3]. Other works related to the decays are given in ref.[4].

**II Two Higgs doublet model**

To begin with, we define the Higgs potential with two Higgs doublets, $\Phi_1$ and $\Phi_2$. We would like to impose the discrete symmetry under $\Phi_2 \rightarrow -\Phi_2$ on the quartic-couplings in the potential to avoid in a natural way the flavor changing neutral current. Then the most general potential becomes

$$V(\Phi_1, \Phi_2) = -\mu_1^2 |\Phi_1|^2 - \mu_2^2 |\Phi_2|^2 - \left( \mu_{12}^2 \Phi_1^\dagger \Phi_2 + \text{h.c.} \right) + \lambda_1 |\Phi_1|^4 + \lambda_2 |\Phi_2|^4 + \lambda_3 |\Phi_1|^2 |\Phi_2|^2 + \lambda_4 \left( \text{Re}\Phi_1^\dagger \Phi_2 \right)^2 + \lambda_5 \left( \text{Im}\Phi_1^\dagger \Phi_2 \right)^2. \quad (\text{II.1})$$

Though the soft breaking term $- \left( \mu_{12}^2 \Phi_1^\dagger \Phi_2 + \cdots \right)$ in eq.(II.1) becomes important in SUSY like models, we here set $\mu_{12}$ for zero because our interests are rather in the strong coupling situation. The effects of $\mu_{12}$ are then suppressed in the heavy mass limit. Therefore, the potential becomes to have seven parameters, $\mu_1, \mu_2, \lambda_1, \sim, \lambda_5$.

Note that the potential (II.1) would have the custodial $SU(2)_V$ symmetry, which is the diagonal part of $SU(2)_L \otimes SU(2)_R$, if $\lambda_5$ would be zero. To see this, it is convenient to rewrite eq.(II.1) in terms of $2 \times 2$ matrices $\Phi_i = (i\tau_2 \Phi_i^r, \Phi_i)$. Then the $\lambda_5$-term in (II.1) becomes

$$\lambda_5 \left\{ \text{tr}(\tau_3 \Phi_1^\dagger \Phi_2) \right\}^2, \quad (\text{II.2})$$
and we can easily see that the term (II.2) breaks $SU(2)_R$ and thus $SU(2)_V$ symmetry explicitly. On the other hand, all the other parts in eq.(II.1) can be rewritten as the combinations of $\text{tr}(\Phi_i^\dagger \Phi_j) (i, j = 1, 2)$, which are clearly custodial $SU(2)_V$ symmetric.

The field configurations in the Higgs doublets are parameterized as

$$\Phi_i = \begin{pmatrix} w_i^+ \\ \frac{1}{\sqrt{2}}(h_i + v_i + iz_i) \end{pmatrix}, (i = 1, 2),$$

where the vacuum expectation values $v_1$ and $v_2$ are combined to give $v = \sqrt{v_1^2 + v_2^2} \sim 246\text{GeV}$. The diagonalization of the mass terms is performed by introducing two kinds of mixing angles $\alpha$ and $\beta$ in the following way;

$$\begin{pmatrix} h_1 \\ h_2 \end{pmatrix} = \begin{pmatrix} \cos \alpha & -\sin \alpha \\ \sin \alpha & \cos \alpha \end{pmatrix} \begin{pmatrix} h \\ H \end{pmatrix},$$

$$\begin{pmatrix} w_1^+ \\ w_2^+ \end{pmatrix} = \begin{pmatrix} \cos \beta & -\sin \beta \\ \sin \beta & \cos \beta \end{pmatrix} \begin{pmatrix} w^+ \\ G^+ \end{pmatrix},$$

$$\begin{pmatrix} z_1 \\ z_2 \end{pmatrix} = \begin{pmatrix} \cos \beta & -\sin \beta \\ \sin \beta & \cos \beta \end{pmatrix} \begin{pmatrix} z \\ A \end{pmatrix}.$$

We set $\tan \beta = v_2/v_1$ as usual, so that fields $w^+, z$ would be Nambu-Goldstone bosons. There are four massive fields, namely, $H, h, G^\pm$ and $A$.

The five quartic-coupling constants in eq. (II.1) are expressed by the masses of these scalar bosons together with the mixing angles;

$$\lambda_1 = \frac{1}{2v^2 \cos^2 \beta}(m_h^2 \cos^2 \alpha + m_H^2 \sin^2 \alpha),$$

$$\lambda_2 = \frac{1}{2v^2 \sin^2 \beta}(m_h^2 \sin^2 \alpha + m_H^2 \cos^2 \alpha),$$

$$\lambda_3 = \frac{\sin 2\alpha}{v^2 \sin 2\beta}(m_h^2 - m_H^2) + \frac{2m_G^2}{v^2},$$

$$\lambda_4 = -\frac{2m_G^2}{v^2},$$

$$\lambda_5 = \frac{2}{v^2}(m_A^2 - m_G^2).$$

Since only the $\lambda_5$-term breaks $SU(2)_V$, eq.(II.3) means that the deviation from the degeneracy between $G^\pm$ and $A$ thus measures the explicit $SU(2)_V$ breaking. The seven
independent parameters (μ₁, μ₂, λ₁, ~, λ₅) in eq.(II.1) are replaced by the four mass parameters (mₕ, m₇, m₆ and m₈), two mixing angles (α and β) and vacuum expectation value v.

III Radiative corrections

Since we assume that all the Higgs masses are much greater than the gauge boson masses, we can make use of the equivalence theorem. The equivalence theorem at loop level is expressed as

\[ T(Z_L(p_1), \cdots, Z_L(p_n), W_L(q_1), \cdots, W_L(q_m); \phi_a) = (C'_Z)^n(C'_W)^m T(iz(p_1), \cdots, iz(p_n), iw(q_1), \cdots, iw(q_m); \phi_a) + O\left(\frac{M_W}{\sqrt{s}}\right), \]

where \( \phi_a \)'s denote the other particles including Higgs bosons, \( \sqrt{s} \) is the typical energy scale of the scattering process and \( C'_Z \) and \( C'_W \) are modification factors to be attached to each external line of \( Z_L \) and \( W_L \)'s respectively. Systematic studies of the general proof of the precise formulation of the equivalence theorem at loop level have been presented by He, Kuang, and Li [7]. According to their work, the modification factors without fermion contributions turn out to be unity if we work in the Landau gauge, on-mass-shell renormalization scheme and heavy-Higgs-mass limit. The important fact is that this statement is kept unchanged even if we work in THDM. The modification factors are, however, to receive additional contributions due to quark loops, which will be discussed in detail in ref.[8]. Thus we have only to calculate the radiative corrections to the processes \( H \to w^+w^- \) and \( zz \) to evaluate those to the processes \( H \to W^+_LW^-_L \) and \( Z_LZ_L \). Though the equivalence theorem does not mention anything about the internal particles, we can neglect all the diagrams with internal \( W^\pm \) and \( Z \) propagators because they are suppressed by \( M_W^2/m_H^2 \). After all, our calculations are reduced simply to those in the Higgs-Goldstone system with top quark.
Now let us prepare the counter-terms for the decay processes. We start from the case of the process $H \to w^+ w^-$. The tree level interaction dictating this process is extracted from the potential (II.1) as

$$L_{Hww} = \frac{m_H^2}{v} \sin(\alpha - \beta) H w^+ w^-.$$  \hspace{1cm} (III.4)

Some of the counter-terms required in the one loop calculation for this process is obtained by varying the parameters in eq.(III.4), namely, by putting $m_H^2 \to m_H^2 - \delta m_H^2, v \to v - \delta v, \alpha \to \alpha - \delta \alpha$, and $\beta \to \beta - \delta \beta$. Others come from the renormalizations of wave-functions and state-mixings between those fields having the same quantum numbers, \textit{i.e.}, by imposing the following replacement upon the bare interaction terms of $H w^+ G^-, HG^+ w^-$ and $hw^+ w^-$ as well as eq.(III.4);

$$\begin{pmatrix} h \\ H \\ w \\ G \\ z \\ A \end{pmatrix} \to \begin{pmatrix} \sqrt{Z_h} \\ \sqrt{Z_{hh}} \\ \sqrt{Z_w} \\ \sqrt{Z_{wG}} \\ \sqrt{Z_z} \\ \sqrt{Z_{zA}} \end{pmatrix} \begin{pmatrix} 1 & -\delta \alpha \\ \delta \alpha & 1 \end{pmatrix} \begin{pmatrix} h \\ H \\ w \\ G \\ z \\ A \end{pmatrix}.$$  \hspace{1cm}

After setting $\sqrt{Z_{hh}} = \sqrt{Z_{hH}}$ and $\sqrt{Z_{wG}} = \sqrt{Z_{Gw}}$, the full counter-term for the process $H \to w^+ w^-$ is obtained as follows;

$$\delta L_{Hww} = \left[ \left( -\frac{\delta m_H^2}{m_H^2} + \frac{\delta v}{v} \right) \frac{m_H^2}{v} \sin(\alpha - \beta) - \frac{m_H^2}{v} (\delta \alpha - \delta \beta) \cos(\alpha - \beta) \right] H w^+ w^-.$$  \hspace{1cm} (III.4)

Similarly, setting $\sqrt{Z_{zz}} = \sqrt{Z_{zA}}$, we also obtain the counter-term for the process, $H \to zz$;

$$\delta L_{Hzz} = \left[ \left( -\frac{\delta m_H^2}{m_H^2} + \frac{\delta v}{v} \right) \frac{m_H^2}{2v} \sin(\alpha - \beta) - \frac{m_H^2}{2v} (\delta \alpha - \delta \beta) \cos(\alpha - \beta) \right] H w^+ w^-.$$  \hspace{1cm} (III.4)
\[
+ \left\{ \left( \sqrt{Z_H} - 1 \right) + (Z_{z} - 1) \right\} \frac{m_H^2}{2v} \sin(\alpha - \beta)
- \left( \sqrt{Z_{hH} - \delta_{\alpha}} \right) \frac{m_h^2}{2v} \cos(\alpha - \beta) \right) \frac{m_H^2 - m_A^2}{v} \cos(\alpha - \beta) \right\} H_{zz}.
\]

We are now full-fledged to perform the one-loop calculations of the amplitudes \(\mathcal{M}_{Hww}(p^2)\) and \(\mathcal{M}_{Hzz}(p^2)\) for the processes \(H \rightarrow w^+w^-\) and \(H \rightarrow zz\). The renormalization is performed in the on-mass shell scheme. Here \(\delta\beta\) is defined by \(zA\) mixing but not by \(wG\) mixing. Details of the calculations will be explained in ref.\[6\]. Finally we arrive at the decay width formula for each process,

\[
\Gamma(H \rightarrow W_L^+W_L^-) = \frac{1}{16\pi m_H} \sqrt{1 - \frac{4M_W^2}{m_H^2}} \left| \mathcal{M}_{Hww}(p^2 = m_H^2) \right|^2 \left| C_{mod}^W \right|^4, \quad (III.5)
\]

\[
\Gamma(H \rightarrow Z_LZ_L) = \frac{1}{32\pi m_H} \sqrt{1 - \frac{4M_Z^2}{m_H^2}} \left| \mathcal{M}_{Hzz}(p^2 = m_H^2) \right|^2 \left| C_{mod}^Z \right|^4. \quad (III.6)
\]

### IV Numerical analysis of the decay widths

In this section, we would like to show some of our numerical analyses for the decay width formulae (III.5) for \(H \rightarrow W_L^+W_L^-\) (seen in part in ref.\[5\]), and (III.6) for \(H \rightarrow Z_LZ_L\). More details of our numerical results will be presented in ref.\[6\].

Some comments on the choice of the parameters are in order. We choose the top-quark mass for 174 GeV and set the mass of the lightest Higgs boson \(m_H\) tentatively for 300 GeV throughout this talk. The mixing angle \(\beta\) is constrained to some extent by the low energy experimental data \[8\], \(\tan\beta\) is not so smaller than unity. We therefore consider either of three cases, \(\tan\beta = 2, 10\) or 20. The mixing angle \(\alpha\) is less bounded phenomenologically \[9\] and so we vary this parameter arbitrarily for theoretical interests. As to the masses of \(h, G^\pm, A\) bosons, these are considerably constrained as a combination with \(m_H\) and mixing angles from the analysis of the \(\rho\) parameter \[10\] and especially \(m_G\) is bounded as \(m_G > 250\text{GeV}\) from the data \[11\]. Here we set \(m_h = 400\text{GeV}\) and the masses of \(G^\pm\) and \(A\) bosons are varied as \(300 < m_G < 900\text{GeV}\) and \(300 < m_A < 1000\text{GeV}\). These parameter
regions are all within the unitarity bounds \[12\]. For the sake of the best illustration of
the features (I) and (II), we mainly show the case \( \sin^2(\alpha - \beta) = 1 \) below. The tree level
evaluation for the decay widths are then calculated as

\[
\Gamma_{\text{tree}}(H \to W_L^+W_L^-) = 7.5\text{GeV}, \quad \Gamma_{\text{tree}}(H \to Z_LZ_L) = 3.5\text{GeV}. \quad (\text{IV.7})
\]

Fig.1 shows the sensitivity of the width formulae (III.5) and (III.6) to \( m_G \) (300 < \( m_G \) < 900GeV) with \( \tan \beta = 2 \) and with \( m_A = 400, 700 \) and 1000GeV (lines a, b and c for \( \Gamma(H \to W_L^+W_L^-) \) and lines \( a', b' \) and \( c' \) for \( \Gamma(H \to Z_LZ_L) \) respectively).

Looking at fig.1, we can see easily that \( \Gamma(H \to Z_LZ_L) \) is quite insensitive to \( m_G \) and
\( m_A \). On the other hand, we can also see that \( \Gamma(H \to W_L^+W_L^-) \) is sensitive to \( m_G \) and
\( m_A \) and that radiative corrections are minimized at \( m_G = m_A \) (recall eqs.(IV.7)). The
difference of the behavior between the decay widths seen in fig.1 will be discussed in the
next section.

Fig.2 shows the mixing angle \( \alpha \) dependence of \( \Gamma(H \to Z_LZ_L) \) for \( m_G = 500, m_A = 600\text{GeV}, \) and \( \tan \beta = 2, 10 \) and 20.
The decay widths as a function of $m_G$. The mixing angles are determined by $\tan \beta = 2$ and $\sin^2(\alpha - \beta) = 1$. The masses of the neutral Higgs bosons are assumed to be $m_H = 300$ GeV and $m_h = 400$ GeV. Lines $a$, $b$, and $c$ represent the behavior of $\Gamma(H \to W^+W^-)$ (eq. (III.5)), while lines $a', b'$ and $c'$ represent that of $\Gamma(H \to Z_LZ_L)$ (eq. (III.6)). The CP-odd Higgs boson mass is taken as $m_A = 400$ GeV ($a$ and $a'$), $m_A = 700$ GeV ($b$ and $b'$) and $m_A = 1$ TeV ($c$ and $c'$), respectively.

V Screening theorem for the vertices

The behavior of the decay widths $\Gamma(H \to W^+_LW^-_L)$ and $\Gamma(H \to Z_LZ_L)$ will be dictated by the leading power contributions with respect to the masses of internal (heavy) scalars $G^\pm, A$ and $h$. We at first would like to discuss these contributions in the amplitudes $\mathcal{M}_{Hww}$ and $\mathcal{M}_{Hzz}$, which are of the form of $M^4/v^3$ (possibly times $\ln M$) on dimensional account ($M$ represents collectively $m_G, m_A$ and/or $m_h$). These contributions are extracted
from the full expression of the amplitudes as

\[ \mathcal{M}_{Hww}(m_H^2) \rightarrow \frac{-1}{(4\pi)^2 v^3} \sin(\alpha - \beta) \]

\[ \times \left\{ (m_A^2 - m_G^2) m_A^2 - m_G^2 \ln \frac{m_A^2}{m_G^2} \right\} \]

\[ + \text{(term from the prescription for } \delta \beta \text{),} \]

(V.8)

\[ \mathcal{M}_{Hzz}(m_H^2) \rightarrow 0, \]

(V.9)

where the second term on RHS in (V.8) has its origin from our prescription scheme for \( \delta \beta \) by the \( zA \)-mixing. This term is extracted from the part which has the factor \( \Pi'_{zA}(0) - \Pi'_{wG}(0) \), where \( \Pi_{zA}(p^2) \) and \( \Pi_{wG}(p^2) \) are two-point functions for \( zA \) and \( wG \) mixing respectively.

The leading contribution (V.8) shows that there are mass-leading contributions in \( \mathcal{M}_{Hww} \) except for the case of \( m_G = m_A \). Recall that \( \Pi_{zA}(p^2) \) would equal \( \Pi_{wG}(p^2) \) for \( m_G = m_A \). This corresponds to the numerical results of \( \Gamma(H \rightarrow W^+_L W^-_L) \), i.e., radiative corrections are sensitive to \( m_G \) and \( m_A \) but insensitive to \( m_h \) and are minimized in the case of \( m_G = m_A \). On the other hand, (V.9) shows that mass-leading contributions in \( \mathcal{M}_{Hzz} \) always cancel out in accordance with the numerical results of \( \Gamma(H \rightarrow Z_L Z_L) \), namely, radiative corrections are always small and insensitive to any of the masses of \( G^\pm, A \) or \( h \).

Now let us consider the reason for these cancellations of the leading-mass contributions in \( \mathcal{M}_{Hww}(p^2) \) for \( m_G = m_A \) and in \( \mathcal{M}_{Hzz} \) at any value of \( m_G \) and \( m_A \). Since the deviation from the mass degeneracy \( m_G = m_A \) measures the custodial symmetry breaking, we may divide each amplitude into two parts as follows,

\[ \mathcal{M}_{Hww} = \mathcal{M}^S + \mathcal{M}^B_{Hww}, \]

(V.10)

\[ \mathcal{M}_{Hzz} = \mathcal{M}^S + \mathcal{M}^B_{Hzz}, \]

(V.11)

where the superscript \( S \) means the custodial \( SU(2)_V \) symmetric part and \( B \) stands for all the other (namely, \( SU(2)_V \) breaking) part. Note that the first terms on RHS in eqs.(V.10) and (V.11) have to be equal to each other owing to the existence of the isospin symmetry.
between \(w^\pm\) and \(z\). The leading contribution (V.8) necessarily comes from \(\mathcal{M}_{Hww}^B\). Both of (V.10) and (V.11) suggests that the mass-leading contributions to \(\mathcal{M}^S\) in eqs. (V.10) and (V.11) always cancel out. We are thus led to the new screening theorem for the vertices which we have presented in Introduction.

The mass-leading contributions seen in (V.8) which have to belong to \(\mathcal{M}_{Hww}^B\) must vanish in the custodial SU(2)\(_V\) symmetric limit. Thus the sensitivity to \(m_G\) and \(m_A\) and the correction-minimization for \(m_G = m_A\) in \(\Gamma(H \rightarrow W_L^+W_L^-)\) are both explained. As to \(\Gamma(H \rightarrow Z_LZ_L)\), the absence of leading-mass contributions, (V.11), which is valid even \(m_G \neq m_A\), indicates that the leading-mass contributions in \(\mathcal{M}_{Hzz}^B\) also cancel out in our scheme (on mass-shell, and definition of \(\delta\beta\) by the \(zA\)-mixing). This non-trivial cancellation in \(\mathcal{M}_{Hzz}^B\) can be proved on the one loop level by making use of renormalizability, absence of the coupling \(zw^+G^-\) and \(zG^+G^-\) in the model and the screening theorem mentioned above. The proof will be presented in ref. [6].

Therefore, we have been able to explain all the characteristics of radiative corrections of the Higgs decay processes \(H \rightarrow W_L^+W_L^-\) and \(Z_LZ_L\) in terms of the screening theorem for vertices.

VI Summary and Discussions

In this talk, we have analyzed the radiative corrections to the decay processes \(H \rightarrow W_L^+W_L^-\) and \(Z_LZ_L\) at one loop level in THDM and have investigated the sensitivity to the masses of the internal scalar bosons. We have found that the radiative corrections to the decay width \(\Gamma(H \rightarrow W_L^+W_L^-)\) are sensitive to the masses of \(G^\pm\) and \(A\) but insensitive to the mass of \(h\) and are minimized for \(m_G = m_A\). On the other hand, it has also been found that the radiative corrections to \(\Gamma(H \rightarrow Z_LZ_L)\) are insensitive to all the masses of internal scalar bosons and are always relatively small.

We have shown that these results are explained completely on the basis of the new
screening theorem for vertices, which applies in the custodial $SU(2)_V$ symmetric limit of the model. This theorem for vertices, though we found it only at one loop level, is also likely to hold to all orders of the perturbation because of the similarity between this theorem and Veltman’s screening theorem for oblique type corrections.

The one-loop insensitivity of $\Gamma(H \to Z_L Z_L)$ to the effect of the internal scalar bosons in our calculation scheme could make this decay width to be a good experimental measuring tool for the mixing angle $\alpha$ (see fig.2). Then the experimental measurement of $\Gamma(H \to W^+_L W^-_L)$ could also provide us with a good measure for the custodial $SU(2)_V$ breaking.

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12
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