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A threshold-improved narrow-width approximation for BSM physics

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Abstract: A modified narrow-width approximation that allows for $\mathcal{O}(\Gamma/M)$-accurate predictions for resonant particle decay with similar intermediate masses is proposed and applied to MSSM processes to demonstrate its importance for searches for particle physics beyond the Standard Model.

Keywords: Beyond Standard Model, NLO Computations, Hadronic Colliders.

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1. Introduction

Theoretical arguments and experimental observations indicate that new particles or interactions play an important role at the TeV scale, which will become directly accessible at the Large Hadron Collider (LHC) – scheduled to start in 2008 – and its planned complement, the International Linear Collider. In the near future we can therefore anticipate groundbreaking discoveries that reveal physics beyond the Standard Model (BSM) and allow to gain insight into the structure of the fundamental theory. Theoretically appealing extensions of the Standard Model (SM) often feature numerous additional interacting heavy particles. Supersymmetric theories [1], for example, are attractive, because they solve the hierarchy problem and allow for the unification of electroweak and strong interactions. The Minimal Supersymmetric Standard Model (MSSM) is one of the best studied candidates for BSM physics. Its phenomenology is characterized by sparticle production and cascade decays, which lead to many-particle final states and scattering amplitudes with complex resonance structure. Cascade decays also occur in other extensions, e.g. in universal extra dimensions models [2].

In order to extract the additional Lagrangian parameters of an extended theory from collider data, theoretical predictions are required that match the experimental accuracies. This can usually only be achieved by taking into account higher order corrections in perturbative calculations. Next-to-leading order calculations for phenomenologically relevant $2 \to n$ processes with $n \gtrsim 4$ are technically very challenging or not yet feasible [3, Sec. 30]. Consequently, production and decay stages are regularly factorized by means of the narrow-width approximation (NWA), which effectively results in on-shell intermediate states.[1] Its main advantage is that sub- and nonresonant as well as nonfactorizable amplitude contributions can be neglected in a theoretically consistent way. Huge calculational simplifications occur already at tree level. For these reasons, the NWA is employed in nearly all studies of BSM physics. Note that it is implicitly applied whenever branching ratios are extracted

[1] The NWA can thus not be applied if on-shell states are kinematically forbidden.

from scattering cross sections. A reliable NWA uncertainty determination is therefore crucial. Given the width Γ and mass M of an unstable particle, the uncertainty of the NWA is commonly estimated as $\mathcal{O}(\Gamma/M)$ for each Breit-Wigner propagator that is integrated out, with Γ/M typically $\lesssim$ 5%. For larger widths nonresonant contributions can no longer be neglected.

Recently, two circumstances have been observed in which the standard NWA is not reliable [4,5]: the first involves decays where a daughter mass m approaches the parent mass M; the second involves the convolution of parton distribution functions with a resonant hard scattering process at center-of-mass energy $\sqrt{\hat{s}}$. In this article we elucidate that both effects arise due to a significant deformation of the Breit-Wigner shape that is caused by threshold factors, and is not restricted to the region where the Breit-Wigner is cut off, i.e. where $M - m$ or $\sqrt{\hat{s}} - M$ is approximately Γ. An essential factor is that the amplitude can contribute additional powers of the threshold factors, which strongly amplifies the effects. For sample applications we then demonstrate that $\mathcal{O}(\Gamma/M)$-accurate predictions can nevertheless be obtained by integrating out the Breit-Wigner in combination with the relevant threshold factors.

2. NWA modifications

To illustrate why the NWA error becomes unexpectedly large for mass configurations in an extended vicinity of kinematical bounds and how it can be modified in such cases, we consider the partial decay rate of a heavy particle A that predominantly decays in two stages via an intermediate resonance C, i.e. $A \xrightarrow{1} B, C$ and $C \xrightarrow{2} D, E$. In terms of the n-body phase space element

$$d\phi(P; p_1, \ldots, p_n) \equiv (2\pi)^4 \delta^{(4)}(P - \sum_{i=1}^{n} p_i) \prod_{i=1}^{n} \frac{d^3 p_i}{(2\pi)^3 2E_i} \tag{2.1}$$

and the matrix element $\mathcal{M}$, the off-shell decay rate is given by

$$\Gamma_{\text{off-shell}} = \frac{1}{2M_A} \int d\phi |\mathcal{M}|^2 \tag{2.2}$$

$$= \frac{1}{2M_A} \int \frac{dp_C^2}{2\pi} D(p_C^2) \int d\phi_1(p_C^2) \int d\phi_2(p_C^2)\ |\mathcal{M}_r(p_C^2)|^2 \ . \tag{2.3}$$

In Eq. (2.3), the phase space factorization

$$d\phi = d\phi_1 \frac{dp_C^2}{2\pi} d\phi_2 \tag{2.4}$$

has been applied, where $d\phi_1(d\phi_2)$ is the 2-body phase space element of the first (second) decay stage. In the rest frame of A,

$$d\phi_1 = \frac{1}{16\pi^2} \frac{|\boldsymbol{p}_C|}{M_A} d\Omega_C \quad \text{with} \quad |\boldsymbol{p}_C| = \frac{M_A}{2} \beta(M_B + \sqrt{p_C^2}, M_A)\, \beta(M_B - \sqrt{p_C^2}, M_A) \ , \tag{2.5}$$

where $\beta(m, M) \equiv \sqrt{1 - m^2/M^2}$. For $d\phi_2$ one finds an analogous expression. In addition to $d\phi$, also $|\mathcal{M}|^2$ has been factorized into the squared propagator denominator

$$D(p_C^2) \equiv \frac{1}{(p_C^2 - M_C^2)^2 + M_C^2 \Gamma_C^2} \tag{2.6}$$

with 4-momentum p_C and $|\mathcal{M}_r|^2$, the residual squared amplitude for the $A \to B, D, E$ decay. Starting in Eq. (2.3), we have indicated the particularly relevant p_C^2-dependence of quantities explicitly, but suppressed the dependence on other kinematical variables. In the limit $\Gamma_C \to 0$, $D(p_C^2)$ is asymptotically equal to $2\pi K \delta(p_C^2 - M_C^2)$ with

$$K = \frac{1}{2 M_C \Gamma_C} = \int_{-\infty}^{\infty} \frac{dq^2}{2\pi} D(q^2) \; . \tag{2.7}$$

This replacement constitutes the standard NWA. Employing it, one obtains

$$\Gamma_{\rm NWA} = \frac{1}{2M_A} K \int d\phi_1(M_C^2) \int d\phi_2(M_C^2) \; |\mathcal{M}_r(M_C^2)|^2 \; . \tag{2.8}$$

It is suggestive to mitigate threshold-induced deviations by absorbing the amplifying factors of β-form that occur in $d\phi_1$ (see Eq. (2.5)) and $d\phi_2$ into K. The absorbed factors have to be normalized to $p_C^2 = M_C^2$, so that they are not taken into account more than once. K is thus replaced by

$$\begin{aligned} \widetilde{K} = & \int_{(M_D+M_E)^2}^{(M_A-M_B)^2} \frac{dp_C^2}{2\pi} \; D(p_C^2) \\ & \times \frac{\beta(M_B + \sqrt{p_C^2}, M_A)\beta(M_B - \sqrt{p_C^2}, M_A)}{\beta(M_B + M_C, M_A)\beta(M_B - M_C, M_A)} \\ & \times \frac{\beta(M_D + M_E, \sqrt{p_C^2})\beta(M_D - M_E, \sqrt{p_C^2})}{\beta(M_D + M_E, M_C)\beta(M_D - M_E, M_C)} \\ & \times \frac{f_{|\mathcal{M}_r|^2}(\sqrt{p_C^2}, M_A, M_B, M_D, M_E)}{f_{|\mathcal{M}_r|^2}(M_C, M_A, M_B, M_D, M_E)} \; . \end{aligned} \tag{2.9}$$

Below we find that additional amplifying factors like $M^2 - m^2 = \beta^2(m, M) M^2$ can arise due to momentum-dependent residual matrix elements. Such factors are included generically in Eq. (2.9) as $f_{|\mathcal{M}_r|^2}$. They are process specific and their powers depend on the spin/polarization of external states.

Note that replacing K with $\widetilde{K}$ does not affect the intrinsic properties of the NWA and that it generalizes to multi-body decays. A closed-form result for $R \equiv \widetilde{K}/K$ can only be given for special cases. In Fig. 1 we show for the ratio R_2 that is obtained by only taking into account the β-factors that arise from $d\phi_2$ and setting the upper integration boundary to infinity the deviation from 1 normalized to Γ_C/M_C. One can see that the largest deviation occurs for $M_D + M_E \to M_C$ and M_E (or M_D) $\to 0$. The sizable effect for $m \equiv M_D \approx M \equiv M_C$ and a small mass M_E, which we set to zero to obtain analytical results, is further amplified if the matrix element of the second decay contributes

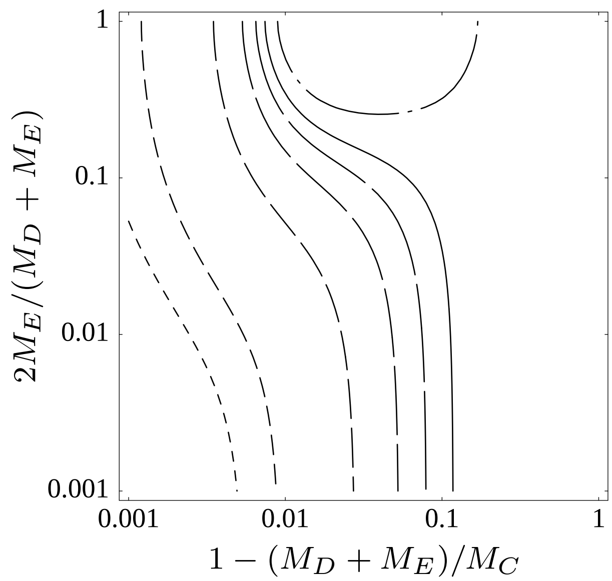

Figure 1: The relative deviation of the modified from the standard NWA factor, normalized to $\Gamma_C/M_C = 1\%$, i.e. $(R_2-1)/(\Gamma_C/M_C)$, is shown as function of the masses of the second decay $C \to D, E$ (see main text for details). Contour lines are shown for the values 0 (solid), -1 (dot-dashed) and $1, 3, 10, 50, 100$ (dashed). The dash length decreases with increasing magnitude.

an additional factor $M^2 - m^2$. This is for example the case if C and E are fermions and D is a scalar (and spin correlations between decay 1 and 2 are neglected) or if C is a scalar and D and E are fermions. For these decay types, strong effects have been observed in Ref. [4] and Ref. [5], respectively. The corresponding $R_2' \equiv \widetilde{K}_2'/K$ is given by

$$\left(\int_{m^2}^{q^2_{\max}} \frac{dq^2}{2\pi} \frac{1}{(q^2-M^2)^2+(M\,\Gamma)^2} \frac{(q^2-m^2)^2/q^2}{(M^2-m^2)^2/M^2}\right) \Bigg/ \left(\int_{-\infty}^{\infty} \frac{dq^2}{2\pi} \frac{1}{(q^2-M^2)^2+(M\,\Gamma)^2}\right) \tag{2.10}$$

with $\Gamma \equiv \Gamma_C$. In Fig. 2 we show the deformation of the Breit-Wigner shape due to the additional threshold factors when m approaches M. After integration, we obtain

$$R_2' = \frac{1}{\pi}\left[\tan^{-1}\frac{\beta^2}{\gamma} + \tan^{-1}\frac{\lambda}{\gamma}\right] + \frac{\gamma}{\pi}\left[\left(\frac{2}{\beta^2}-1\right)\ln\frac{\lambda}{\beta^2} + \left(\frac{1}{\beta^2}-1\right)^2 \ln\frac{q^2_{\max}}{m^2}\right] \tag{2.11}$$

with $\gamma \equiv \Gamma/M$, $\beta \equiv \beta(m, M)$ and $\lambda \equiv q^2_{\max}/M^2 - 1$ when factors of $1+\gamma^2$ are approximated by 1. (This approximation does not produce a visible difference in Fig. 3.) The result confirms that away from threshold, where $\beta \approx 1$, one obtains $R_2' \approx 1$ with $\gamma \ll \beta^2, \lambda$. When approaching the threshold, i.e. $\beta \to 0$, the divergence of the second, formally γ-suppressed term overcompensates the decrease of the first term. In Fig. 3 we show the β

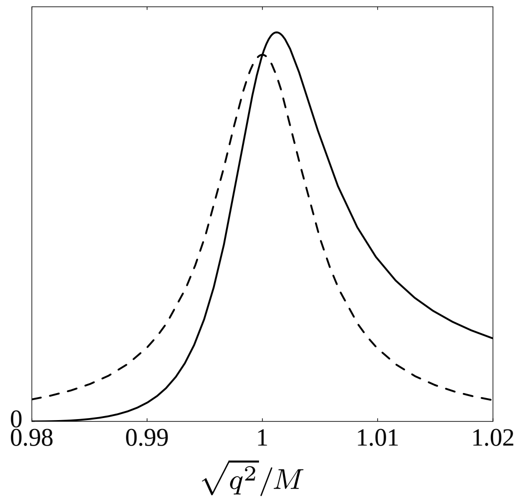

Figure 2: The Breit-Wigner shape deformation is displayed that is caused by threshold factors when a decay daughter mass m approaches the parent mass M. More specifically, the integrand of the numerator (solid) and denominator (dashed) of Eq. (2.10) are shown in unspecified normalization as functions of the invariant mass $\sqrt{q^2}$. $\Gamma/M = 1\%$ and $m = M - 2\Gamma$.

dependence of R_2' for typical values of $\sqrt{q^2_{\max}}/M$. The deviation from the standard NWA clearly exceeds $\mathcal{O}(\Gamma/M)$ already for threshold masses m that are still significantly below the resonant region roughly bounded by $M \pm \Gamma$. We note that if $\sqrt{q^2_{\max}}/M - 1 \lesssim \gamma$

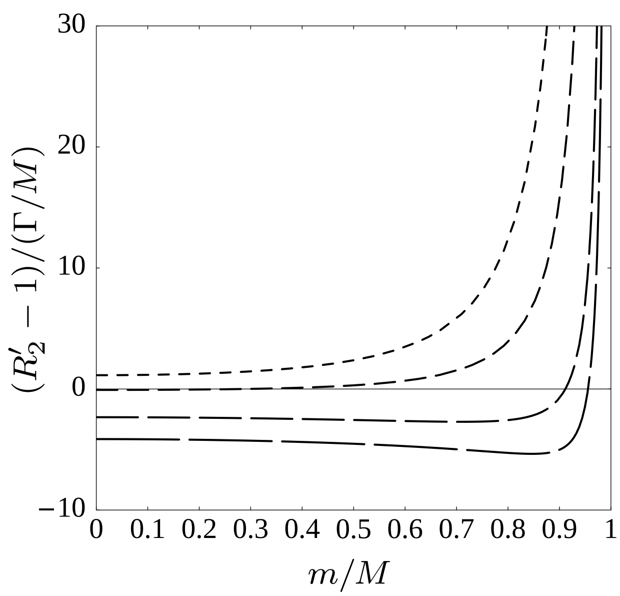

Figure 3: The decay mass dependence of the relative deviation of the modified from the standard NWA factor is shown in units of the conventionally expected uncertainty $\Gamma/M = 1\%$ for $\sqrt{q^2_{\max}}/M \in \{1.05, 1.1, 2, 10\}$. The dash length decreases with increasing $\sqrt{q^2_{\max}}$.

the threshold amplification is confined to the resonant region. However, in this case the arctan terms are no longer approximately $\pi/2$, which results in a much larger than expected uncertainty of the standard NWA for arbitrary values of m. If $q^2_{\rm max} \gg M^2$ the contribution from the region $q^2 \approx q^2_{\rm max}$ to R'_2 is enhanced by the factor $(q^2 - m^2)^2/(q^2m^2) \approx q^2/m^2$ and the production cross section's suppression close to threshold can become important.

3. Applications

In this section we demonstrate that the NWA modification proposed in Sec. 2 allows to reduce the uncertainty to the conventional expectation for mass configurations in the vicinity of kinematical bounds.

As a first application we study the threshold-improved approximation for scalar scattering and decay. More specifically, we study the processes displayed in Fig. 4. For this type of process large standard NWA deviations have been observed in Refs. [4, 5] when $m_d \approx M$ or the center of mass energy $\sqrt{s} \approx M$. We assess the quality of the modified

Figure 4: Process 1 (left) with scalar scattering and decay and process 2 (right) with decay into non-scalar particles (fermions). Lines without labels correspond to massless particles.

NWA by comparing the off-shell cross section to the cross section $\sigma_{\rm INWA}$ calculated in NWA with $\widetilde{K}$ of Eq. (2.9), where $M_A = \sqrt{s}, M_C = M, \Gamma_C = \Gamma, M_D = m_d$ and $M_B = M_E = 0$. The deviation is measured in units of Γ/M, which we set to 0.01, using

$$R_{\rm INWA} \equiv \left(\frac{\sigma_{\text{off-shell}}}{\sigma_{\rm INWA}} - 1\right) \Big/ \frac{\Gamma}{M} \ . \tag{3.1}$$

We start by neglecting matrix element effects and thus set $f_{|\mathcal{M}_r|^2} = 1$. For process 1 with $m_p \sim M$ we find satisfactory NWA uncertainty reduction. For instance for $m_p = 1.1M$, $R_{\rm INWA} \lesssim 3$ as long as $(\sqrt{s} - M)/M \gtrsim 10^{-5}$. For $m_p \ll M$, however, large deviations occur in a significant parameter space region, in particular for $m_d \approx M$. For process 2 large deviations remain, independent of the value of m_p. The Breit-Wigner deformation arises apparently not just from threshold-type phase space element factors. To achieve a satisfactory NWA uncertainty reduction it is in general essential to take into account factors originating from the matrix element that distort the Breit-Wigner shape. We separate them into production and decay-related factors:

$$f_{|\mathcal{M}_r|^2}(M, \sqrt{s}, 0, m_d, 0, m_p) = f_p(M, \sqrt{s}, m_p)\ f_d(M, m_d) \ . \tag{3.2}$$

For process 1, the decay matrix element is a coupling constant and we thus have $f_d = 1$. For process 2, however, we have

$$f_d(M, m_d) = \frac{|\mathcal{M}_d|^2}{m_d^2} = \frac{M^2 - m_d^2}{m_d^2} = \beta^2(m_d, M)\frac{M^2}{m_d^2} \;, \tag{3.3}$$

where we have divided by m_d^2 to obtain a dimensionless quantity and expressed the result in terms of threshold β-factors that also appear in the decay phase space element (see Eq. (2.5)) in order to show the deviation-amplifying effect of the decay matrix element by contributing an additional power to the Breit-Wigner deforming factor in $d\phi_d$. When taking decay matrix element effects into account, i.e. employing $\widetilde{K}$ with $f_{|\mathcal{M}_r|^2} = f_d(M, m_d)$, NWA deviations are mitigated to $\mathcal{O}(\Gamma/M)$, except for the region $\sqrt{s} \lesssim 1.5M$, where the t-channel production matrix element causes significant Breit-Wigner deformations. These production effects can be remedied with

$$f_p(M, \sqrt{s}, m_p) = \beta^{-2}\left(\sqrt{|M^2 - m_p^2|}, \sqrt{s}\right) \;. \tag{3.4}$$

We now extend our analysis to more complex processes and study the NWA uncertainty reduction at hadron colliders for sparticle production and decay in the MSSM. In Ref. [4] the standard NWA accuracy was studied for the process $u\bar{d} \to (\tilde{g} \to \widetilde{s}_{L,R}\bar{s})\,\widetilde{\chi}_1^+$. Here, the resonant particle, i.e. the gluino, is produced in a t-channel process with either $\tilde{d}_L$ or $\tilde{u}_L$ exchange. For this process a variation of the $\widetilde{s}_L$ mass between 0 and the gluino mass revealed unexpectedly large NWA deviations for squark masses that are larger than $0.8M_{\tilde{g}}$. The slope of the increasing deviation when the squark mass approaches the gluino mass is qualitatively very similar to the slope displayed for R_2' in Fig. 3. Since R_2' does not take into account the t-channel production effects, we conclude that they do not significantly alter the dominant decay effects. We have confirmed that in this region the uncertainty of the NWA is reduced to $\mathcal{O}(\Gamma/M)$ if $\widetilde{K}_2'$ of Eq. (2.10) is used with $M = M_{\tilde{g}}$, $\Gamma = \Gamma_{\tilde{g}}$, $m = M_{\widetilde{s}}$ and $\sqrt{q^2_{\text{max}}} = 1.4M_{\tilde{g}}$ (matched at the squark mass value where $\sigma_{\text{off-shell}} = \sigma_{\text{NWA}}$). We note that for $\widetilde{s}_L$ masses below $0.8M_{\tilde{g}}$ the NWA overestimates the off-shell cross section by up to about 20%. This deviation is, however, consistent with an expected uncertainty of $\mathcal{O}(\Gamma/M)$, since in this region the gluino width increases to about 10% of its mass. Ref. [4] also illustrates deviations for the $\widetilde{s}_L$-$\widetilde{s}_R$ decay asymmetry, which are consistent with NWA corrections of $\mathcal{O}(\Gamma/M)$.

As a last application we consider cascade decays, which are the natural testing ground for Eq. (2.9). More specifically, we study $\tilde{g}\,\tilde{u}_L$ production at the LHC, i.e. in proton-proton collisions at 14 TeV, with the subsequent cascade decay $\tilde{g} \to \widetilde{s}_L\bar{s}$ and $\widetilde{s}_L \to \widetilde{\chi}_1^- c$ at the SPS1a' benchmark point [6] in the MSSM parameter space. Phenomenologically, to consider a squark decay into the LSP candidate $\widetilde{\chi}_1^0$ would be more natural, but the resulting complete Feynman amplitude features a complicated resonance structure whose study we leave to future work. Even for the gluino decay chain considered here, interference arises from $\tilde{g} \to (\widetilde{c}_L^* \to \widetilde{\chi}_1^-\bar{s})c$. Its effect is, however, small. We confirmed that omitting it does not affect our $\mathcal{O}(\Gamma/M)$ accuracy goal. In this article we focus on the resonant $\widetilde{s}_L$ state (with $M = 570$ GeV and $\Gamma = 5.4$ GeV at SPS1a') and the NWA accuracy relative to Γ/M

that is obtained using $\widetilde{K}$ with

$$f_{|\mathcal{M}_r|^2}(M_{\tilde{s}}, M_{\tilde{g}}, 0, M_{\widetilde{\chi}}, 0) = \frac{M_{\tilde{g}}^2 - M_{\tilde{s}}^2}{M_{\tilde{g}}^2}\,\frac{M_{\tilde{s}}^2 - M_{\widetilde{\chi}}^2}{M_{\widetilde{\chi}}^2} = \beta^2(M_{\tilde{s}}, M_{\tilde{g}})\,\beta^2(M_{\widetilde{\chi}}, M_{\tilde{s}})\,\frac{M_{\tilde{s}}^2}{M_{\widetilde{\chi}}^2} \qquad (3.5)$$

versus K when the strange squark mass approaches either the gluino or chargino mass of 607 and 184 GeV, respectively.[2] Results calculated with MadEvent [7] and Sherpa [8] using CTEQ6L1 parton distribution functions [9] and spectra and decay widths obtained with SPheno [10] and SDECAY [11] are displayed in Fig. 5. The Monte Carlo integration error

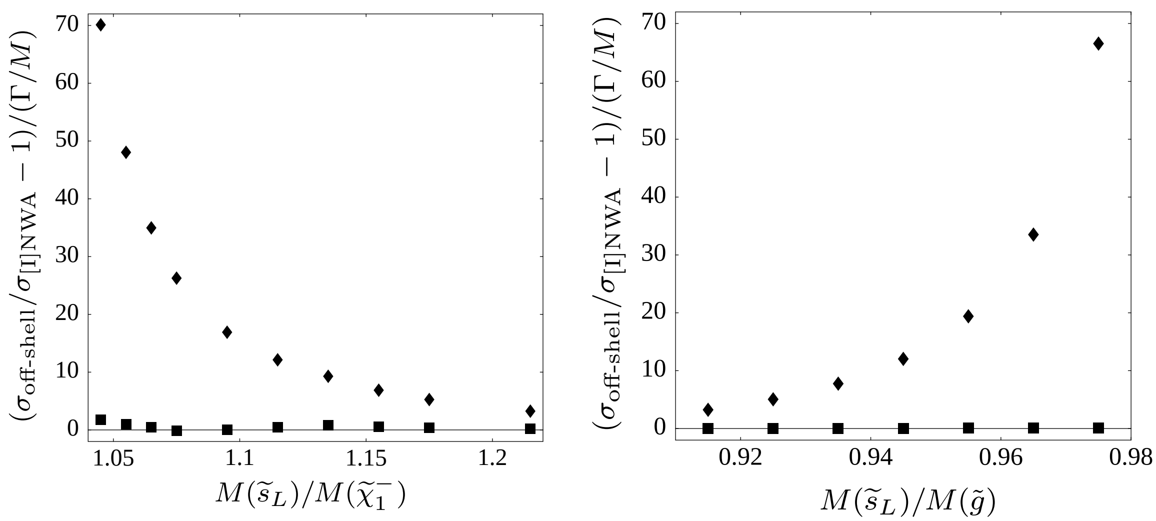

Figure 5: The accuracy of the NWA cross section normalized to the conventionally expected uncertainty is shown for $\tilde{g}\,\tilde{u}_L$ production at the LHC followed by the cascade decay $\tilde{g} \to \tilde{s}_L\bar{s}$ and $\tilde{s}_L \to \widetilde{\chi}_1^- c$ in the MSSM at SPS1a' for a variable strange squark mass that approaches the chargino mass (left) and the gluino mass (right). Results are displayed for the standard NWA (diamonds) and the improved NWA (INWA) of Eq. (2.9) (boxes). $\Gamma(\tilde{s}_L)/M(\tilde{s}_L)$ ranges from 0.03% to 0.16% (left) and is approximately 0.9% (right).

is 0.1%. Both figures show that the modified NWA reduces the sizable deviations that occur in standard NWA as a daughter or parent mass is approached to the conventional uncertainty estimate. A multiple, overlapping application of Eq. (2.9) that would also include gluino production and chargino decay effects could be envisioned, but is beyond the scope of this work.

4. Conclusions

For configurations with kinematical bounds in the vicinity of resonances phase space suppression via β-factors can significantly distort the resonance Breit-Wigner, thus effecting an unexpectedly large NWA error. For affected configurations we proposed a modification

[2]The chargino is treated as stable and the gluino in standard NWA with spin correlations.

of the standard NWA that allows to take this kinematical phase space suppression into account and thus to reduce the approximation uncertainty to the inverse of the generic resonant enhancement M/Γ. For supersymmetric extensions of the SM we have demonstrated this uncertainty reduction for similar masses in processes with single particle or cascade decay. If applied in phenomenological studies and data analysis with tools like Fittino [12], SFITTER [13] or MARMOSET [14], the method would contribute to an accurate determination of BSM model parameters and thus to establishing supersymmetry or other key properties of the fundamental theory.

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References

[1] H. P. Nilles, Phys. Rept. **110** (1984) 1; H. E. Haber and G. L. Kane, Phys. Rept. **117** (1985) 75.

[2] C. Macesanu, Int. J. Mod. Phys. A **21** (2006) 2259 [arXiv:hep-ph/0510418].

[3] C. Buttar *et al.*, arXiv:hep-ph/0604120.

[4] D. Berdine, N. Kauer and D. Rainwater, Phys. Rev. Lett. **99** (2007) 111601 [arXiv:hep-ph/0703058].

[5] N. Kauer, Phys. Lett. B **649** (2007) 413 [arXiv:hep-ph/0703077].

[6] J. A. Aguilar-Saavedra *et al.*, Eur. Phys. J. C **46** (2006) 43 [arXiv:hep-ph/0511344]; B. C. Allanach *et al.*, Eur. Phys. J. C **25** (2002) 113 [arXiv:hep-ph/0202233].

[7] F. Maltoni and T. Stelzer, JHEP **0302** (2003) 027 [arXiv:hep-ph/0208156]; T. Stelzer and W. F. Long, Comput. Phys. Commun. **81** (1994) 357 [arXiv:hep-ph/9401258]; J. Alwall *et al.*, JHEP **0709** (2007) 028 [arXiv:0706.2334 [hep-ph]].

[8] T. Gleisberg, S. Hoche, F. Krauss, A. Schalicke, S. Schumann and J. C. Winter, JHEP **0402** (2004) 056 [arXiv:hep-ph/0311263]; F. Krauss, R. Kuhn and G. Soff, JHEP **0202** (2002) 044 [arXiv:hep-ph/0109036]; K. Hagiwara *et al.*, Phys. Rev. D **73** (2006) 055005 [arXiv:hep-ph/0512260].

[9] J. Pumplin, D. R. Stump, J. Huston, H. L. Lai, P. Nadolsky and W. K. Tung, JHEP **0207** (2002) 012 [arXiv:hep-ph/0201195].

[10] W. Porod, Comput. Phys. Commun. **153** (2003) 275 [arXiv:hep-ph/0301101].

[11] M. Muhlleitner, A. Djouadi and Y. Mambrini, Comput. Phys. Commun. **168** (2005) 46 [arXiv:hep-ph/0311167].

[12] P. Bechtle, K. Desch and P. Wienemann, Comput. Phys. Commun. **174** (2006) 47 [arXiv:hep-ph/0412012].

[13] R. Lafaye, T. Plehn and D. Zerwas, arXiv:hep-ph/0404282.

[14] N. Arkani-Hamed, P. Schuster, N. Toro, J. Thaler, L. T. Wang, B. Knuteson and S. Mrenna, arXiv:hep-ph/0703088.