Hunting physics beyond the standard model with unusual $W^{\pm}$ and Z decays

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Nonstandard on-shell decays of $W^{\pm}$ and Z bosons are possible within the framework of extended supersymmetric models, i.e., with singlet states and/or new couplings compared to the minimal supersymmetric standard model. These modes are typically encountered in regions of the parameter space with light singlet-like scalars, pseudoscalars, and neutralinos. In this letter we emphasize how these states can lead to novel signals at colliders from Z- or $W^{\pm}$-boson decays with prompt or displaced multileptons/tau jets/jets/photons in the final states. These new modes would give distinct evidence of new physics even when direct searches remain unsuccessful. We discuss the possibilities of probing these new signals using the existing LHC run-I data set. We also address the same in the context of the LHC run-II, as well as for the future colliders. We exemplify our observations with the "μ from ν" supersymmetric standard model, where three generations of right-handed neutrino superfields are used to solve shortcomings of the minimal supersymmetric standard model. We also extend our discussion for other variants of supersymmetric models that can accommodate similar signatures.

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I. INTRODUCTION

Physics beyond the standard model (SM) remains necessary even after the long-awaited discovery of the Higgs boson [1]. A much anticipated but hitherto unseen excess over the SM thus makes it rather essential and timely to explore other methods, e.g., precision measurements of the SM observables, where evidence of new physics may remain hidden. In this article we present an analysis of this kind in the context of extended supersymmetric (SUSY) models, which can lead to new two-body decays of $W^{\pm}$ and Z bosons for certain regions of the *enlarged* parameter space. Concerning a Z boson, decaying to a scalar and a pseudoscalar, we search for final states with a combination of four prompt leptons ($\ell = e, \mu$)/τ jets (from hadronically decaying τ)/jets/photons. We look for similar but displaced final states with missing transverse energy ($\not{E}_{\mathrm{T}}$) when a Z boson decays into a pair of neutralinos. For the $W^{\pm}$ boson, when it decays into a charged lepton and a neutralino, we look for final states with two displaced leptons/τ jets/jets/photons + $\not{E}_{\mathrm{T}}$ together with a prompt lepton/τ jet.

Earlier analyses of unusual $W^{\pm}$ decays [2] in the minimal supersymmetric standard model (MSSM) are already excluded by the chargino mass bound [3]. Unusual Z decays for models with an extended Higgs sector with/without SUSY have also been studied, both theoretically [4] and experimentally [5]. These decays lead to either multiparticles or missing transverse momentum/energy signatures at colliders. Most of these scenarios—e.g., a light ($< M_Z/2$) doublet-like pseudoscalar as discussed in the third and sixth papers of Ref. [4]—hardly survive with the current experimental constraints. However, a study of all new $W^{\pm}$ and Z two-body decays through singlet-like states in light of the Higgs boson discovery [1] is missing to date. In this article we aim to carry out this analysis using as a case study the "μ from ν" supersymmetric standard model ($\mu\nu$SSM) [6,7], the *simplest* variant of the MSSM to house these signals, apart from housing nonzero neutrino masses and mixing [6–11] and offering a solution to the μ problem [12] of the MSSM. Nonetheless, we also extend our discussion to other model variants. Note that the said light singlet-like states also contribute to the invisible/nonstandard decay branching fractions (BRs) for a Higgs-like scalar [13,14].

II. THE $\mu\nu$SSM

In the $\mu\nu$SSM, three families of right-handed neutrino superfields ($\hat{\nu}_i^c$) are instrumental in offering a solution to the μ problem [12] of the MSSM, and concurrently in housing the observed pattern of neutrino masses and mixing [6–11]. The superpotential is given by

$$W = \epsilon_{ab}(Y_{u_{ij}} \hat{H}_u^b \hat{Q}_i^a \hat{u}_j^c + Y_{d_{ij}} \hat{H}_d^a \hat{Q}_i^b \hat{d}_j^c + Y_{e_{ij}} \hat{H}_d^a \hat{L}_i^b \hat{e}_j^c + Y_{\nu_{ij}} \hat{H}_u^b \hat{L}_i^a \hat{\nu}_j^c - \lambda_i \hat{\nu}_i^c \hat{H}_d^a \hat{H}_u^b) + \frac{1}{3}\kappa_{ijk} \hat{\nu}_i^c \hat{\nu}_j^c \hat{\nu}_k^c. \tag{1}$$

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The last three terms in Eq. (1) break R parity ($\not{R}_p$) explicitly. After the spontaneous breaking of the electroweak symmetry, the neutral scalars develop vacuum expectation values as $\langle H_d^0 \rangle = v_d$, $\langle H_u^0 \rangle = v_u$, $\langle \tilde{\nu}_i \rangle = \nu_i$, and $\langle \tilde{\nu}_i^c \rangle = \nu_i^c$; thus, they generate the effective bilinear terms $\varepsilon_i \hat{L}_i \hat{H}_u$, $\mu \hat{H}_d \hat{H}_u$ and Majorana mass terms $M_{ij} \hat{\nu}_i^c \hat{\nu}_j^c$, with $\varepsilon_i \equiv Y_{\nu_{ij}} \nu_j^c$, $\mu \equiv \lambda_i \nu_i^c$, and $M_{ij} \equiv 2\kappa_{ijk} \nu_k^c$.

The enlarged field content and the $\not{R}_p$ in the $\mu\nu$SSM result in eight CP-even (S_α^0) and seven CP-odd (P_α^0) neutral scalar, seven charged scalar ($S_\alpha^\pm$), ten neutralino ($\tilde{\chi}_\alpha^0$), and five chargino ($\tilde{\chi}_\alpha^\pm$) states [6–8,11]. The three lightest neutralinos and charginos, denoted as $\tilde{\chi}_i^0$ and $\tilde{\chi}_i^\pm$, coincide with the left-handed neutrinos and the charged leptons.

In the $\mu\nu$SSM, light (in order to trigger new Z, $W^\pm$, and Higgs decays, i.e., $\lesssim M_Z/2$ [9,15–18]) right-handed sneutrino ($\tilde{\nu}^c$) and right-handed neutrino (ν_R)-like states, in the bottom of the mass spectrum, are possible with suitable choices of λ_i, κ_{ijk}, $\tan\beta (= \frac{v_u}{v_d})$, ν_i^c, and the soft SUSY-breaking parameters [6,7] $A_{\kappa_{ijk}}$ and A_{λ_i}. The parameters λ_i and A_{λ_i} control the doublet impurity and hence affect the lightness of these states [16–18]. The parameters κ_{ijk}, $A_{\kappa_{ijk}}$, and ν_i^c determine their mass scales [17,18]. A small doublet component, i.e., a small λ_i, together with small $\tan\beta$ values makes it easier (see, e.g., Ref. [19]) for these states to evade a class of collider [20–22] and low-energy constraints [23–25] (see Ref. [18] for details). Here, S_i^0, P_i^0, and $\tilde{\chi}_{i+3}^0$ are used to denote $\tilde{\nu}^c$-like scalars, pseudoscalars, and ν_R-like neutralinos in the mass eigenbasis, respectively. With this notation, S_4^0 is the SM-like Higgs boson [7,15–18]. The effects of S_i^0, P_i^0, and $\tilde{\chi}_{i+3}^0$ states in S_4^0 decays have already been addressed in the $\mu\nu$SSM before [15,16] and after [17,18] the Higgs discovery [1].

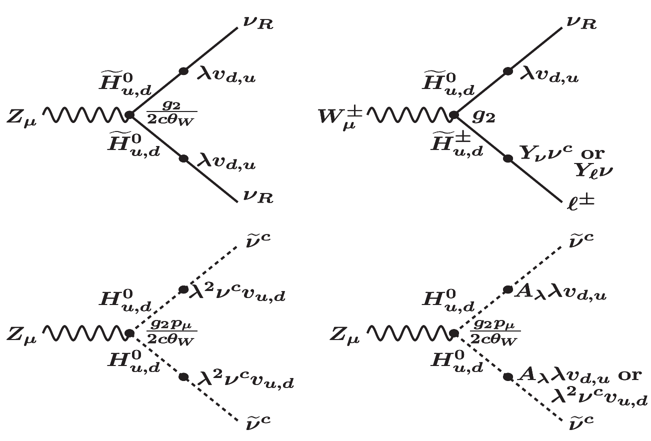

FIG. 1. Leading on-shell $W^\pm$ and Z decays through ν_R and $\tilde{\nu}^c$, where g_2 is the SU(2) gauge coupling, θ_W is the Weinberg angle, c $\equiv$ cos, and p_μ is the momentum factor for the $ZS_i^0 P_j^0$ vertices. Family indices and extra factors coming from the field decomposition (as mentioned in the text) are not explicitly shown.

TABLE I. Table showing all possible final states from non-standard Z and $W^\pm$ decays along with their respective origins.

Z decay	$W^\pm$ decay
$2x^D 2\bar{x}^D + \not{E}_T$(via $\tilde{\chi}_{i+3}^0 \tilde{\chi}_{j+3}^0$)	ℓ^P/τ-jet$^P + x^D \bar{x}^D$
$2x^P 2\bar{x}^P$(via $S_i^0 P_j^0$)	$+\not{E}_T$(via $\tilde{\chi}_i^\pm \tilde{\chi}_{j+3}^0$)

III. NEW Z,$W^\pm$ DECAYS IN THE $\mu\nu$SSM

New on-shell decays, like $Z \to S_i^0 P_j^0$, $\tilde{\chi}_{i+3}^0 \tilde{\chi}_{j+3}^0$ and $W^\pm \to \tilde{\chi}_i^\pm \tilde{\chi}_{j+3}^0$, as already stated, naturally open up for $m_{S_i^0, P_i^0, \tilde{\chi}_{i+3}^0} \lesssim M_Z/2$. These decays are schematically shown in Fig. 1 using the flavor basis. Depending on $m_{S_i^0, P_i^0}$ [18], successive S_i^0, P_i^0 decays into a pair of leptons/τ jets/jets/photons give prompt multiparticle final states for the $Z \to S_i^0 P_j^0$ processes. A light neutralino (with a mass < 20 GeV) in SUSY models with singlets can decay through lighter [9]/slightly heavier [17] $S_i^0/P_i^0 + \tilde{\chi}_j^0$ modes. These new modes dominate for $m_{\tilde{\chi}^0} < 40$ GeV and can give displaced decays within the tracker [9,16,17]. The subsequent S_i^0/P_i^0 decay gives a combination of four displaced but detectable leptons/τ jets/jets/photons [17,18] for the $Z \to \tilde{\chi}_{i+3}^0 \tilde{\chi}_{j+3}^0$ modes, with some $\not{E}_T$ from neutrinos and possible mis-measurements. Note that the $S_4^0 \to S_i^0 S_j^0$, $P_i^0 P_j^0$, and $\tilde{\chi}_{i+3}^0 \tilde{\chi}_{j+3}^0$ decays [18] produce final states that are similar to those of $Z \to S_i^0 P_j^0$, $\tilde{\chi}_{i+3}^0 \tilde{\chi}_{j+3}^0$ processes. The $W^\pm \to \tilde{\chi}_i^\pm \tilde{\chi}_{j+3}^0$ decays, in the same way, lead to final states with one prompt lepton/τ jet + two displaced leptons/τ jets/jets/photons + $\not{E}_T$. Possible final states are shown below in tabular format (see Table I) with $x =$ lepton/τ jet/jet/photon. The superscripts P and D are used to denote prompt and displaced nature, respectively.

These new decays are constrained by the measured decay widths, i.e., $\Gamma_Z = 2.4952 \pm 0.0023$ GeV, $\Gamma_{W^\pm} = 2.085 \pm 0.042$ GeV, and $\Gamma_Z^{\text{inv}} = 0.499 \pm 0.0015$ GeV [3] if $\tilde{\chi}_{i+3}^0$ decays invisibly, i.e., outside the detector. From the latest theoretical calculation, including higher-order contributions, one gets $\Gamma_Z^{\text{theo}} = 2.49424 \pm 0.0005$ GeV and $\Gamma_Z^{\text{inv(theo)}} = 0.50147 \pm 0.000014$ GeV at the 1σ level [26]. So one can compute $\delta\Gamma_Z \equiv \Gamma_Z - \Gamma_Z^{\text{theo}} = 1.0 \pm 2.4$ MeV and, similarly, $\delta\Gamma_Z^{\text{inv}} = -2.5 \pm 1.5$ MeV, where we have added the theoretical and experimental 1σ errors in quadrature. It is thus apparent that at the 1σ level[1] one has a freedom of about 3.4 MeV to accommodate the new physics contributions in Z decays over the SM prediction. Concerning new invisible Z decays beyond the SM, one is forced to consider a 2σ variation in order to accommodate a

[1]We restrict ourselves within 68% C.L. of errors (see 11th paper of Ref. [4]), which gives more stringent constraints.

positive new physics effect with a freedom of about 0.5 MeV. A similar analysis for $W^\pm$ with $\Gamma^{\rm theo}_{W^\pm} = 2.0932 \pm 0.0022$ GeV [27] gives $\delta\Gamma_{W^\pm} = -8 \pm 42$ MeV. Thus, any new contribution must give a decay width less than about 34 MeV at the 1σ level. Moreover, from the SM one gets $\mathrm{BR}(Z \to 4\ell^P) = 4.2^{+0.9}_{-0.8} \times 10^{-6}$ [3,28] and $\mathrm{BR}(Z \to 4b^P) = 3.6^{+1.3}_{-1.3} \times 10^{-4}$ [3,29]. The latest $\mathrm{BR}(Z \to 4\ell^P)$ is estimated as $3.2^{+0.28}_{-0.28} \times 10^{-6}$ [30].

Regarding final states, $S^0_i, P^0_i \to b\bar{b}$ remains generic for $2m_b \lesssim m_{S^0_i,P^0_i} \lesssim M_Z/2$ while taus dominate for $2m_\tau \lesssim m_{S^0_i,P^0_i} \lesssim 2m_b$ [16–18]. Light jets ($\not\ni b$) or leptons normally emerge for $m_{S^0_i,P^0_i} \lesssim 2m_\tau$ with a moderate parameter tuning, especially for λ_i, κ_{ijk}, ν^c_i, and $A_{\kappa ijk}$ [18]. A multiphoton signal through the $S^0_i/P^0_i \to \gamma\gamma$ process (that is *unique* in nature) requires large parameter tuning to get a statistically significant result. Note that a low mass of the mother particle, e.g., $M_Z \approx 91$ GeV or $M_W \approx 80$ GeV, typically gives rise to soft leptons/τ jets/jets/photons, in the final states [17,18], i.e., with low transverse momentum, that suffer from poor detection efficiency [31].

Concerning the lightest neutralino $\tilde{\chi}^0_4$, a Higgsino- or wino-like nature is forbidden for $2m_{\tilde{\chi}^0_4} \lesssim M_Z$ from the lighter chargino mass bound [3]. The latter demands the minimum of $(\mu, M_2) \gtrsim 100$ GeV, where M_2 is the SU(2) gaugino soft mass. Further, the tree-level Z-Higgsino-Higgsino interaction constrains the amount of Higgsino impurity in $\tilde{\chi}^0_4$ from the measured Γ_Z [3]. So, depending on the relative orders of $2\kappa_{ijk}\nu^c_k$ and M_1 [U(1) gaugino soft mass], a light $\tilde{\chi}^0_4$ is either bino- or right-handed neutrino-like, or a bino-right-handed neutrino mixed state. A light bino-like $\tilde{\chi}^0_4$ needs a low M_1 and hence a relaxation of the gaugino mass unification at the high scale when considered together with the LHC bound on gluino mass [32]. Besides, as lighter S^0_i, P^0_i states are not assured for a bino-like $\tilde{\chi}^0_4$ [18], typical decay occurs beyond the tracker and often outside the detector [9,16]. Hence we do not consider this possibility. One can also consider a bino-like $\tilde{\chi}^0_7$, with $2m_{\tilde{\chi}^0_7} \lesssim M_Z$, that decays to right-handed neutrino-like $\tilde{\chi}^0_{i+3} + S^0_j/P^0_j$, followed by $\tilde{\chi}^0_{i+3} \to \tilde{\chi}^0_j + S^0_k/P^0_k$, and finally $S^0_i/P^0_i \to$ a pair of leptons/τ jets/jets/photons. In this scenario, in spite of the large particle multiplicity of the final state, most of these leptons/τ jets/jets/photons remain undetected due to their soft nature [18], and thus this is not studied here.

IV. NEW $Z, W^\pm$ DECAYS IN OTHER MODELS

Unusual Z decays with prompt final states also appear in R_p-conserving models with singlets, e.g., the next-to-MSSM (NMSSM) (see fifth paper of Ref. [19] for a review). The presence of $3\hat{\nu}^c_i$ in the $\mu\nu$SSM, which emerges naturally from the family symmetry, can produce different peaks in the invariant mass ($m_{\rm inv}/\mathrm{M_{T2}}$) [33] distributions for the two-leptons/τ jets/jets/photons systems. In the NMSSM, however, one expects a single peak with one singlet superfield. With dedicated search strategies, e.g., better detection efficiency for a soft lepton/τ jet/jet/photon, etc., and higher statistics one could hope to probe (and hence discriminate) these scenarios in the coming years. The model of Ref. [34] in the context of the studied signals is similar to the $\mu\nu$SSM, although additional constraints can appear for the former from dark matter searches. Note that in SUSY models with singlets, $S^0_4 \to S^0_i S^0_j$, $P^0_i P^0_j$, and $\tilde{\chi}^0_{i+3}\tilde{\chi}^0_{j+3}$ processes [16,18,35] can mimic new Z decays, but they give a different $m_{\rm inv}/M_{\rm T2}$ peak for the four-particle system around $m_{S^0_4}$.

Displaced $W^\pm$ and Z decays are also possible in the MSSM with $\not{R}_p$ [36], sometimes with a richer topology, e.g., two-leptons/τ jets + 4 jets + $\not{E}_{\rm T}$ with $\lambda'_{ijk}\hat{L}_i\hat{Q}_j\hat{d}^c_k$ couplings. For the bilinear $\not{R}_p$ model, a light neutralino with a mass < 20 GeV normally decays outside the detector [9]. The MSSM with trilinear $\not{R}_p$ couplings [$\lambda_{ijk}\hat{L}_i\hat{L}_j\hat{e}^c_k$ for $\lambda_{ijk} \sim \mathcal{O}(10^{-3})$], however, can produce a decay length within 1 cm–3 m [37] for a neutralino in the same mass range. Nevertheless, with different intermediate states it is possible to identify (and hence discriminate) these signals by constructing a set of kinematical variables. However, one needs a good reconstruction efficiency for the displaced and normally soft lepton/τ jet/jet/photon to probe these signals.

Displaced Z decays in the NMSSM can appear for fine-tuned λ values [38], when a pair of next-to-lightest SUSY particles (NLSPs) are produced in the Z decay. These NLSPs, when decay further into a lightest supersymmetric particle and a scalar/pseudoscalar, followed by the scalar/pseudoscalar decays into two leptons/τ jets/jets/photons produce identical displaced final states. The corresponding decay length, however, is never in the range of 1 cm–3 m (second paper of Ref. [38]).

Look-alike displaced states also appear for the model of Ref. [34] even with an R_p-conserving vacuum, when a Z decays into two singlino-like NLSPs, followed by an NLSP $\to \tilde{\nu}^c$-like lightest supersymmetric particle and a right-handed neutrino process. This right-handed neutrino then decays further to a scalar/pseudoscalar and a left-handed neutrino, giving an identical signal. However, in this case $\not{E}_{\rm T}$ could be larger and the scenario is constrained from dark matter searches.

These observations verify the $\mu\nu$SSM as the *minimal extension* beyond the MSSM to house these distinctive collider signatures together with the correct neutrino physics, as well as offering a solution to the μ problem.

V. THE BACKGROUNDS

Leading SM backgrounds for $Z \to S^0_i P^0_j$ processes come through Drell-Yan, $W^\pm W^\mp$, $W^\pm Z$, ZZ/γ, $b\bar{b}$, dileptonically decaying $t\bar{t}$, $W^{\pm}/Z+$ jets and $Z\to 4l$, 4 jets, $2l2$ jets, where l represents a charged lepton. A faithful reconstruction of $m_{\rm inv}/{\rm M}_{\rm T2}$ for the $4l/4$ jets$/2l2$ jets as well as $2l/2$ jets systems, with a proper experimental setup (as mentioned earlier), can disentangle these signals. The latter is crucial to isolate the studied signals from $Z\to 4l$, 4 jets, $2l2$ jets backgrounds [39], where the $m_{\rm inv}/M_{\rm T2}$ distribution for four-particle systems also peaks around M_Z. A recent analysis [30] only considered e and μ, while in the $\mu\nu$SSM a $\tau(b)$ rich signal is generic for $2m_\tau \lesssim m_{S_i^0,P_i^0} \lesssim 2m_b (m_{S_i^0,P_i^0} \gtrsim 2m_b)$. Displaced Z, $W^{\pm}$ decays are exempted from the SM backgrounds.

VI. ESTIMATING NEW BRs AND DECAY WIDTHS

Turning back to the $\mu\nu$SSM, we aim to estimate BRs for nonstandard $W^{\pm}, Z$ decays analytically, and hence we stick to the flavor basis for an easy interpretation. Z decays—which are suppressed by small $Y_{\nu_{ij}}/\nu_i$, as required for a TeV-scale seesaw [6–11]—are not shown in Fig. 1. The smallness of $Y_{\nu_{ij}}$ and ν_i assure that BRs for the tree-level flavor-violating $Z\to \tilde{\chi}_i^{\pm}\tilde{\chi}_j^{\mp}$ processes are well below the respective SM limits. Note that the same logic also predicts very small BRs for flavor-violating Higgs decays into a pair of leptons. Hence, it remains difficult to accommodate the recent excess in the $\mu\tau$ final state from Higgs decays, as reported by the CMS Collaboration [40]. For $W^{\pm}$, on the contrary, processes involving $Y_{\nu_{ij}}/\nu_i$ are the leading nonstandard decay modes, and hence they are always severely suppressed by orders of magnitude compared to the unusual Z decays.

We start our discussion with the complete expressions for the new $W^{\pm}$ and Z decay widths, i.e., $\Gamma(W^{\pm}\to\tilde{\chi}_i^{\pm}\tilde{\chi}_{j+3}^0)$, $\Gamma(Z\to\tilde{\chi}_{i+3}^0\tilde{\chi}_{j+3}^0)$, and $\Gamma(Z\to S_i^0P_j^0)$. These are written as follows:

$$\begin{aligned}
\Gamma(W^{\pm}\to\tilde{\chi}_i^{\pm}\tilde{\chi}_{j+3}^0) &= \frac{g_2^2}{48\pi M_W^5}\left[\left(M_W^2-m^2_{\tilde{\chi}_i^{\pm}}-m^2_{\tilde{\chi}_{j+3}^0}\right)^2-4m^2_{\tilde{\chi}_i^{\pm}}m^2_{\tilde{\chi}_{j+3}^0}\right]^{\frac{1}{2}}\times\Big[\Big((|O^{cnw}_{Lij+3}|^2+|O^{cnw}_{Rij+3}|^2)\\
&\quad\times\left(2M_W^4-M_W^2\left(m^2_{\tilde{\chi}_i^{\pm}}+m^2_{\tilde{\chi}_{j+3}^0}\right)-\left(m^2_{\tilde{\chi}_i^{\pm}}-m^2_{\tilde{\chi}_{j+3}^0}\right)^2\right)+12\Re\left(O^{cnw^*}_{Lij+3}O^{cnw}_{Rij+3}\right)m_{\tilde{\chi}_i^{\pm}}m_{\tilde{\chi}_{j+3}^0}M_W^2\Big],\\
\Gamma(Z\to\tilde{\chi}_{i+3}^0\tilde{\chi}_{j+3}^0) &= \frac{g_2^2}{48\pi M_Z^5}\left[\left(M_Z^2-m^2_{\tilde{\chi}_{i+3}^0}-m^2_{\tilde{\chi}_{j+3}^0}\right)^2-4m^2_{\tilde{\chi}_{i+3}^0}m^2_{\tilde{\chi}_{j+3}^0}\right]^{\frac{1}{2}}\times\Big[\Big((|O^{nnz}_{Li+3j+3}|^2+|O^{nnz}_{Ri+3j+3}|^2)\\
&\quad\times\left(2M_Z^4-M_Z^2\left(m^2_{\tilde{\chi}_{i+3}^0}+m^2_{\tilde{\chi}_{j+3}^0}\right)-\left(m^2_{\tilde{\chi}_{i+3}^0}-m^2_{\tilde{\chi}_{j+3}^0}\right)^2\right)+12\Re\left(O^{nnz^*}_{Li+3j+3}O^{nnz}_{Ri+3j+3}\right)m_{\tilde{\chi}_{i+3}^0}m_{\tilde{\chi}_{j+3}^0}M_Z^2\Big],\\
\Gamma(Z\to S_i^0P_j^0) &= \frac{g_2^2|O^{spz}_{ij}|^2}{192\pi\cos^2\theta_{\rm W}M_Z^5}\left[\left(M_Z^2-m^2_{S_i^0}-m^2_{P_j^0}\right)^2-4m^2_{S_i^0}m^2_{P_j^0}\right]^{\frac{1}{2}}\times\left[M_Z^4-2M_Z^2\left(m^2_{S_i^0}+m^2_{P_j^0}\right)+\left(m^2_{S_i^0}-m^2_{P_j^0}\right)^2\right].
\end{aligned}\tag{2}$$

The couplings $O^{nnz}_{L(R)}$, $O^{cnw}_{L(R)}$ are given in Ref. [11], and $O^{spz}_{ij}=\{R^{S^0}_{i1}R^{P^0}_{j1}-R^{S^0}_{i2}R^{P^0}_{j2}+\sum_{p=1}^{3}R^{S^0}_{i,p+5}R^{P^0}_{j,p+5}\}$, where $R^{S^0}_{ab}$ $(R^{P^0}_{ab})$ denotes the amount of the bth flavor state in the ath scalar (pseudoscalar) mass eigenstate after rotating away the Goldstone boson. We will use these full formulas in a forthcoming publication [41] for a complete numerical analysis of the unusual $W^{\pm}, Z$ decays over the different regions of the parameter space. Thus, in this article we will use a set of approximate formulas for various decay widths to explore the behavior of these new decays analytically.

The relative strength of the three $Z\to S_i^0P_j^0$ processes in Fig. 1, assuming $\nu^c\approx A_\lambda$, goes as $\lambda^2{:}1{:}\lambda$. Clearly, the process $\propto A_\lambda^2$ dominates over the rest when λ is small ($\lesssim 0.1$), unless $\nu^c\gtrsim 10A_\lambda$. A small λ also ensures singlet purity for the light S_i^0, P_i^0, and $\tilde{\chi}_{i+3}^0$ states [7,9,15–18]. Note that the decay BR of a light S_i^0, P_i^0, or $\tilde{\chi}_{i+3}^0$ state into a specific mode can be tuned to 100% depending on $m_{S_i^0,P_i^0}$ [16,18]. In this derivation and for the subsequent analysis we assume, for simplicity, universal λ_i, $A_{\lambda i}$, ν_i^c, i.e., $\mu=3\lambda\nu^c$ and $(\kappa,A_\kappa)_{ijk}=(\kappa,A_\kappa)_i\delta_{ij}\delta_{jk}$, with $\kappa_i\approx\kappa_j$ [16,17]. Further, we take $(Y_\nu, A_\nu[6,7])_{ij}=(Y_\nu,A_\nu)_i\delta_{ij}$. These approximate formulas are also valid in the mass basis, provided that the doublet contamination in S_i^0 and P_i^0 states is negligible.

$\Gamma(Z\to\tilde{\chi}_{i+3}^0\tilde{\chi}_{j+3}^0)$ and $\Gamma(Z\to S_i^0P_j^0)$, following their orders in Fig. 1, are estimated as

$$\begin{aligned}
\Gamma_{Z\tilde{\chi}^0\tilde{\chi}^0} &\approx \frac{6g_2^2\lambda^4v_u^4M_Z^2}{2^6c^2_{\theta_W}\mu^4}\mathcal{P}, \qquad \Gamma^1_{ZS^0P^0}\approx\frac{6g_2^2\lambda^4\mu^4v_u^4M_Z^2}{2^63^4c^2_{\theta_W}\tilde{m}^8}\mathcal{P},\\
\Gamma^2_{ZS^0P^0} &\approx \frac{6g_2^2A_\lambda^4\lambda^4v_u^4M_Z^2}{4c^2_{\theta_W}\tilde{m}^8}\mathcal{P}, \qquad \Gamma^3_{ZS^0P^0}\approx\frac{6g_2^2A_\lambda^2\lambda^4\mu^2v_u^4M_Z^2}{2^43^2c^2_{\theta_W}\tilde{m}^8}\mathcal{P},
\end{aligned}\tag{3}$$

where $\mathcal{P} = \frac{1}{16\pi m_A}\sqrt{\{1-(\frac{m_B^2}{m_A^2}+\frac{m_C^2}{m_A^2})\}^2 - 4\frac{m_B^2}{m_A^2}\frac{m_C^2}{m_A^2}}$ is the phase-space factor for a $A \to BC$ process. A factor of "6" in the numerator comes after summing over all possible i and j values without double counting.[2] We have also introduced a factor (not shown in Fig. 1) coming from the field decomposition as discussed in Ref. [18] following the expressions given in Ref. [7]. Those factors are 2^4, 2^4, 1, and 2^2, respectively, for the four decay widths shown in Eq. (3). The generic mass scale for the intermediate Higgsino (Higgs) is denoted by $\mu(\tilde{m})$. For Z decays at rest, $p_\mu p^\mu = M_Z^2$.

In order to estimate the maximum new Γ_Z we have neglected contributions $\propto v_d$ compared to v_u for $\tan\beta > 1$, and thus we have used $v_u \approx v = 174$ GeV. Now with $g_2 = 0.652$, $\lambda = 0.1$, $c_W^2 = 0.769$, and $M_Z = 91.187$ GeV [3], and assuming $\mu, A_\lambda, \tilde{m} \approx \mathcal{O}(v)$, the total new Γ_Z (from three leading contributions, i.e., $\Gamma_{Z\tilde{\chi}^0\tilde{\chi}^0} + \Gamma^2_{ZS^0P^0} + \Gamma^3_{ZS^0P^0}$) from Eq. (3) is evaluated as ≈ 0.16 MeV. This number is smaller than 3.4 MeV, as estimated in Sec. III. The total new Γ_Z in this region of parameter space for $\lambda \gtrsim 0.21$ becomes larger than 3.4 MeV and hence experimentally disfavored. Further, we also note that the lowest possible value for the μ parameter is $\approx$100 GeV, as required from the lighter chargino mass bound [3]. In this region of parameter space, i.e., when $\mu \approx 100$ GeV while $\tilde{m}, A_\lambda \approx \mathcal{O}(v)$, the total new Γ_Z is larger than 3.4 MeV for $\lambda \gtrsim 0.19$. In this estimation we assume $m_{S_i^0} \sim m_{P_i^0} \sim m_{\tilde{\chi}^0_{i+3}}$ and consider $\mathcal{P} \approx 1/16\pi M_Z$. This approximation remains valid as long as the S_i^0, P_i^0, and $\tilde{\chi}^0_{i+3}$ states are much lighter than M_Z. For heavier S_i^0, P_i^0, and $\tilde{\chi}^0_{i+3}$ states, the new Γ_Z reduces due to a suppression from the phase-space factor. We note in passing that in this analysis we focus only on "visible" $\tilde{\chi}^0_{i+3}$ decays. The corner of the parameter space (in the case of genuinely invisible Z decays) with $\mu \lesssim 130$ GeV lowers the upper limit of λ (e.g., $\lesssim 0.16$ using $\mu = 100$ GeV) from a freedom of 0.5 MeV in the measurement of the invisible Z-decay width. For a larger value of the μ parameter, $\Gamma_{Z\tilde{\chi}^0\tilde{\chi}^0}$ appears to be quite suppressed compared to $\Gamma^2_{ZS^0P^0}$ [see Eq. (3)], and hence the constraint on the total Z-decay width is reached faster than that of the invisible Z decay. The latter observation reverses for $\tilde{m} \gtrsim 1.3\mu$ with $A_\lambda \approx \tilde{m}$ irrespective of the magnitude of A_λ, μ, and $\tilde{m}$ with respect to the scale of v.

One can, nevertheless, use $A_\lambda, \tilde{m}, \mu > v$ to suppress the new Γ_Z as well as the new $\Gamma_Z^{\rm inv}$ for any λ values. This way, larger λ values can survive the experimental constraint on Γ_Z, e.g., $A_\lambda, \tilde{m}, \mu = 2v$ gives a new $\Gamma_Z < 3.4$ MeV even if $\lambda = 0.4$. Such large λ values,[3] however, spoil the singlet purity and lightness of S_i^0, P_i^0, and $\tilde{\chi}^0_{i+3}$ states [18].

So we conclude that, for $\lambda \lesssim 0.1$ and $\mathcal{O}(1\text{TeV}) \gtrsim A_\lambda, \tilde{m}, \mu \gtrsim \mathcal{O}(v)$, one gets $\mathcal{O}(10^{-8}) \lesssim \text{BR}(Z \to S_i^0 P_j^0, \tilde{\chi}^0_{i+3}\tilde{\chi}^0_{j+3}) \lesssim \mathcal{O}(10^{-4})$. Further, $\text{BR}(S_i^0/P_i^0 \to x^P \bar{x}^P)$ and $\text{BR}(\tilde{\chi}^0_{i+3} \to x^D \bar{x}^D + \not{E}_{\rm T}) \approx 1$ imply that $\text{BR}(Z \to 2b^P 2\bar{b}^P)$ and $\text{BR}(Z \to 2b^D 2\bar{b}^D + \not{E}_{\rm T})$ also vary in the *same* range. Thus, the new decay BR for the $Z \to 2b^P 2\bar{b}^P$ process (i.e., when $2m_b \lesssim m_{S_i^0, P_i^0} \lesssim M_Z/2$) remains *below* the SM measured value [3,29], while $Z \to 2\ell^P 2\bar{\ell}^P$ (with $m_{S_i^0, P_i^0} \lesssim 2m_\tau$) remains comparable (at the 2σ level) or below the concerned SM limit [30] for $A_\lambda, \tilde{m}, \mu \geq 356$ GeV. For the region of parameter space with $\mu \approx 100$ GeV and $A_\lambda, \tilde{m} \approx \mathcal{O}(v)$, one needs to consider $\lambda \lesssim 0.044$ to respect the SM measurement of $\text{BR}(Z \to 2\ell^P 2\bar{\ell}^P)$, concerning the 2σ variation around the central value [30]. These issues will be addressed in detail in Ref. [41]. The other way of getting $Z \to 2\ell^P 2\bar{\ell}^P$—i.e., through leptonic tau decays following $Z \to 2\tau 2\bar{\tau}$ and taking the leptonic τ decay $\text{BR} \approx 0.35$ [3]—gives $\mathcal{O}(10^{-10}) \lesssim \text{BR}(Z \to 2\ell^P 2\bar{\ell}^P) \lesssim \mathcal{O}(10^{-6})$, which is less than or comparable to the measured SM value [3,28,30]. Note that $Z \to 2\ell 2\tau, 4\tau$ decays are experimentally unconstrained to date.

Let us finally remark that $\lambda \sim 0.1$ and $\mu \gtrsim 100$ GeV imply $\nu^c \approx 333$ GeV. Hence, $0.1 \text{ GeV} \lesssim |m_{\tilde{\chi}^0_{i+3}}| \lesssim 45$ GeV predicts $0.00015 \lesssim \kappa_i \lesssim 0.07$. Similarly, approximate formulas for $m_{P_i^0}$ give $0.07 \text{ GeV} \lesssim |A_{\kappa i}| \lesssim 29$ GeV. The relative sign difference between κ and A_κ predicts $0.08 \text{ GeV} \lesssim m_{S_i^0} \lesssim 37$ GeV, which does not introduce a large error in the $m_{S_i^0} \sim m_{P_i^0}$ assumption. Here we have used $m^2_{S_i^0} \approx m^2_{\tilde{\chi}^0_{i+3}} + (\kappa A_\kappa)_i \nu^c$, $m^2_{P_i^0} \approx -3(\kappa A_\kappa)_i \nu^c$, and $m_{\tilde{\chi}^0_{i+3}} \approx 2\kappa_i \nu^c$ in the limit of vanishingly small λ.

Concerning $W^\pm$, leading decay widths are approximately given by

$$\Gamma^1_{W^\pm\tilde{\chi}^\mp\tilde{\chi}^0} \approx \frac{9g_2^2 Y_\nu^2 v_u^2 M_W}{144\pi\mu^2} \times \left(1 - \frac{m^2_{\tilde{\chi}^0_{j+3}}}{M_W^2}\right),$$

$$\Gamma^2_{W^\pm\tilde{\chi}^\mp\tilde{\chi}^0} \approx \frac{9g_2^2 \lambda^2 Y_l^2 v_u^2 \nu^2 M_W}{16\pi\mu^4} \times \left(1 - \frac{m^2_{\tilde{\chi}^0_{j+3}}}{M_W^2}\right), \qquad (4)$$

where we have skipped terms like $m^2_{\tilde{\chi}_i^\pm}/M_W^2$ due to their smallness and the factor "9" in the numerator appears after

[2]One can write a set of similar formulas for "n" families of right-handed neutrino superfields when the numerator looks like $(\frac{n+1}{2n})\ \boldsymbol{\lambda}^4$. Here $\boldsymbol{\lambda}^2 \equiv \sum_i \lambda_i^2 = n\lambda^2$, assuming universal λ_i. The quantity $\boldsymbol{\lambda}^2$ is bounded from above from the requirement of maintaining the perturbative nature of λ_i parameters up to some energy scale, e.g., $\lesssim (0.7)^2$ when the scale lies at 10^{16} GeV [7]. Hence, adding more and more singlets, i.e., $n \to \infty$ does not imply a blow-up behavior for the new Z-decay widths.

[3]Note that by assuming the perturbative nature of λ_i parameters up to the grand unified theory scale, i.e., 10^{16} GeV, one gets $\boldsymbol{\lambda}^2 \lesssim (0.7)^2$ [7]. This predicts a maximum for $\lambda \approx 0.4$ assuming universal λ_i. Higher λ_i values require a lower scale up to which the concerned parameters respect their perturbative nature. [7].

summing over all possible combinations. Now, as before, $\mu \approx \mathcal{O}(v)$, $\lambda = 0.1$, and $M_W = 80.385$ GeV [3] with a light $\tilde{\chi}^0_{j+3}$ and a maximum of $(Y_\nu, \nu) \sim (10^{-6}, 10^{-4}$ GeV) for a TeV-scale seesaw [6–11] give a maximum total new $\Gamma_{W^\pm} \sim \mathcal{O}(10^{-9}$ MeV) $\ll 34$ MeV, as evaluated in Sec. III. This new $\Gamma_{W^\pm}$ decreases further with larger μ values. Note that unlike the new Z decays, we do not consider extra factors coming from the field decomposition (see Ref. [18]) which would produce further suppression. Here we assume $Y_\ell \sim \mathcal{O}(1)$ (which holds true for $\tan\beta \gg 1$) *only* for a maximum estimate. Hence, together with $\mathrm{BR}(\tilde{\chi}^0_{i+3} \to x^D \bar{x}^D + \not{E}_{\mathrm{T}}) \approx 1$, one gets $\mathrm{BR}(W^\pm \to \ell^{\pm P}/\tau^{\pm P} + x^D \bar{x}^D + \not{E}_{\mathrm{T}}) \sim 3 \times 10^{-13}$ as a maximum. This BR can at most reach $\mathcal{O}(10^{-12})$ when $\mu \approx 100$ GeV.

VII. PROBING NEW $Z, W^\pm$ DECAYS AT COLLIDERS

Production cross sections for $W^\pm$ and Z at the LHC run-I (center-of-mass energy $\mathrm{E_{CM}} = 8$ TeV) are estimated as $\sigma_{\mathrm{prod}} \approx 8.6 \times 10^4$ and 2.5×10^4 pb, respectively. PYTHIA (version 6.409) [42] has been used for this purpose (verified with MADGRAPH5 version 1.4.2 [43]), with leading-order CTEQ6L1 parton distribution functions [44] having initial- and final-state radiations and multiple interactions switched on. For run-II, i.e., $\mathrm{E_{CM}} = 13$ and 14 TeV, σ_{prod} scales by a factor of ~ 2.

The huge BR suppression for new $W^\pm$ decays predicts ≈ 0.15 events [$\mathcal{O}(1)$ events with $\mathrm{BR} \sim \mathcal{O}(10^{-12})$ at the parton level], taking $\mathrm{E_{CM}} = 14$ TeV and an integrated luminosity $\mathcal{L} = 3000\ \mathrm{fb}^{-1}$. Clearly, even without further reduction from the efficiency issues, these practically background-free signals are rather difficult to detect with both current and upcoming collider searches (e.g., *MegaW* and *OkuW* modes of the Linear Collider [45] and triple large electron–positron collider (TLEP) [46] with about 2×10^6 and 7×10^8 $W^\pm$ bosons/year, respectively). Detecting these signals may appear feasible by using the techniques from the flavor sector, especially when $\mathrm{BR}(B^0_s \to \mu^+\mu^-)$ and $\mathrm{BR}(\mu \to e\gamma)$ have already been probed up to $\mathcal{O}(10^{-10})$ [47] and $\mathcal{O}(10^{-13})$ [48], respectively. The latter would reach $\mathcal{O}(10^{-19})$ in the near future [49].

On the other hand, unusual Z decays (both prompt and displaced) for the LHC run-I with $\mathcal{L} = 25\ \mathrm{fb}^{-1}$ gives about 62 500 parton-level events with the maximum BR estimate. Note that, as already stated, not all of the daughter particles are hard, i.e., with high transverse momentum. Hence, this number will reduce further considering practical issues like b-tagging (τ-tagging) efficiency [31], faking, etc. For displaced decays, the reconstruction efficiency of the displaced vertices diminishes this number even further. Assuming an optimal scenario with a detection efficiency of 50% (25%) for the two leading (subleading) daughters, experimentally one can detect only $\approx 1.5\%$ of the parton-level events, i.e., about 938 surviving events. Repeating the same exercise for the LHC run-II with $\mathrm{E_{CM}} = 13$ TeV and $\mathcal{L} = 100\ \mathrm{fb}^{-1}$, one gets about 7500 surviving events out of 5×10^5 parton-level events. Making a similar analysis for $Z \to 4b^P$ in the SM with a BR of $3.6^{+1.3}_{-1.3} \times 10^{-4}$ gives 3375^{+1219}_{-1219} and 27000^{+9750}_{-9750} events for run-I and run-II, respectively.

It is now apparent that the final number of events for these spectacular signatures remains below that coming from the errors of the SM measurement. Thus, unless one adopts dedicated experimental searches (as already mentioned) to perform a faithful construction of the $2b$ jets m_{inv} ($\mathrm{M_{T2}}$ for τ jets), which is supposed to peak around $m_{S^0_i, P^0_i}$, it remains hardly possible to isolate these new signals from the SM backgrounds. Consequently, the region of parameter space with $2m_b \lesssim m_{S^0_i, P^0_i} \lesssim M_Z/2$ remains unconstrained from the existing experimental results. A similar logic remains applicable for the $Z \to 4\ell^P$ process, through leptonic τ decays following $Z \to 4\tau^P$. Here, for run-I, and $\mathcal{L} = 20.3\ \mathrm{fb}^{-1}$ as studied by the ATLAS Collaboration [30], one gets a maximum of 761 events at the parton level, including e, μ flavors, while the SM number is 1624^{+142}_{-142} with the latest $\mathrm{BR}(Z \to 4\ell^P)$ [30]. So, unless $\mu, A_\lambda, \tilde{m} \approx \mathcal{O}(v)$, these novel signals normally lie beneath their SM counterparts, and hence the region of parameter space giving $2m_\tau \lesssim m_{S^0_i, P^0_i} \lesssim 2m_b$ remains exempted from the experimental constraints. The existing results for $\mathrm{BR}(Z \to 4\ell^P)$ measurements in the SM [3,28,30], however, put constraints on the corner of parameter space with $m_{S^0_i, P^0_i} \lesssim 2m_\tau$, unless $A_\lambda, \tilde{m}, \mu > v$ or $\lambda \ll 0.1$. The $m_{S^0_i, P^0_i} \lesssim 2m_\tau$ region, however, is already challenged by a class of flavor observables [23].

Note that, concerning statistics, the upcoming colliders (*GigaZ* and *TeraZ* modes of the Linear Collider [45] and TLEP [46] with about 2×10^9 and 7×10^{11} Z bosons/year, respectively) are comparable to the LHC run-II, e.g., with *TeraZ* one expects about 10^4 to 10^8 novel Z-decay events when the BR varies from 10^{-8} to 10^{-4}. Nonetheless, the unprecedented accuracy of these new colliders would proficiently constrain the concerned region of parameter space giving rare signals, e.g., in TLEP the error in Γ_Z would reduce to < 10 keV [46].

VIII. CONCLUSIONS

In conclusion, we have presented a complete analytical study of all possible two-body nonstandard $W^\pm$ and Z decays in the context of SUSY theories. The detection of these decays at the LHC or in a future collider (of course, with evolved search criteria) would provide an unambiguous sign of new physics beyond the SM. The presence of these signals will also offer an experimental test for the $\mu\nu$SSM, the *simplest* extension of the MSSM to accommodate these unique signatures, apart from solving shortcomings of the MSSM. We also show that other variants of SUSY models can be investigated in a similar way. After the latest and improved measurement of $Z \to 4\ell^P$ processes by the ATLAS Collaboration, we look forward to further analysis in this direction, both from the CMS and ATLAS Collaborations, involving b jets or τ jets, coming from both prompt and displaced origins. Missing evidence of these modes, in the presence of a proper experimental setup, would constrain the feasibility of light S_i^0, P_i^0, and $\tilde{\chi}_{i+3}^0$ states. Even the existing bound on the branching fractions can disfavor certain regions of parameter space, e.g., λ around 0.1 with $\mu \sim 100$ GeV. We further observed that the region of parameter space with $\mu, A_\lambda, \tilde{m} \approx \mathcal{O}(v)$ and $m_{S_i^0, P_i^0} \lesssim 2m_\tau$ is largely excluded from $Z \to 4\ell^P$ searches in the SM unless one considers $\lambda \ll 0.1$.

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