New supersymmetric Janus solutions from $N = 4$ gauged supergravity

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We study $N = 4$ gauged supergravity with an $SO(4) \times SO(4)$ gauge group in the presence of symplectic deformations and find new classes of Janus solutions preserving $N = 1$ and $N = 2$ supersymmetries. The $N = 2$ solutions preserve $SO(2) \times SO(2) \times SO(2) \times SO(2)$ symmetry and interpolate between $N = 4$ supersymmetric AdS_4 vacua with $SO(4) \times SO(4)$ symmetry. These correspond holographically to $N = (2, 0)$ two-dimensional conformal defects within the dual $N = 4$ Chern-Simons-Matter (CSM) theories with $SO(4) \times SO(4)$ symmetry. The $N = 1$ solutions contain two families of Janus configurations, one interpolating between $N = 4$ AdS_4 vacua with $SO(4) \times SO(4)$ symmetry and the other interpolating between $N = 4$ AdS_4 vacua with $SO(3) \times SO(3) \times SO(3)$ symmetry. These describe, respectively, $N = (1, 0)$ conformal defects in $N = 4$ CSM theories with $SO(4) \times SO(4)$ and $SO(3) \times SO(3) \times SO(3)$ symmetries. The latter give the first example of Janus solutions involving nontrivial AdS_4 vacua in addition to the trivial $SO(4) \times SO(4)$ critical point at the origin of the scalar manifold within the framework of $N = 4$ gauged supergravity.

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I. INTRODUCTION

Janus configurations are solutions of gauged supergravity theories in the form of anti–de Sitter (AdS)-sliced (curved) domain walls interpolating between AdS vacua. According to the AdS/CFT correspondence [1–3], these solutions holographically describe conformal interfaces or defects within the dual conformal field theories [4]; see also [5–8]. These defects break the conformal symmetry of the bulk superconformal field theory (SCFT) down to that on the codimension-1 defects by some position-dependent operators; see [9,10] for recent results. For almost 20 years since the first Janus solution of [4], a large number of Janus solutions have been found in gauged supergravities in various space-time dimensions with different numbers of supersymmetries; see [11–37] for an incomplete list.

In this paper, we are interested in supersymmetric Janus solutions from symplectically deformed $N = 4$ gauged supergravity with an $SO(4) \times SO(4)$ gauge group. The $N = 4$ gauged supergravity coupled to n vector multiplets was constructed in the embedding tensor formalism in [38] (see [39–41] for earlier constructions), and possible symplectic deformations were considered in [42], extending the construction of ω-deformed $SO(8)$ maximal gauged supergravity [43–46] to lower numbers of supersymmetry. As pointed out in [42], for $N = 4$ gauged supergravity with an $SO(4)\times SO(4) \sim SO(3)\times SO(3)\times SO(3)\times SO(3)$ gauge group, there can be four deformation parameters or electric-magnetic phases for each $SO(3)$ factor. The first two $SO(3)$ factors are embedded in $SO(6)_R \sim SU(4)_R$ R symmetry of $N = 4$ supersymmetry. One of the phases for this $SO(3) \times SO(3)$ can be set to zero by $SL(2, \mathbb{R})$ transformations of the global symmetry $SL(2, \mathbb{R}) \times SO(6, n)$, while the other gives equivalent gauged supergravities for any nonvanishing values and can be set to $\frac{\pi}{2}$. The phases of the remaining two $SO(3)$ factors embedded in the $SO(n)$ symmetry of the matter vector multiplets are independent deformation parameters, in constrast to a single phase ω of the maximal $SO(8)$ gauged supergravity. The vacuum structure of the symplectically deformed $SO(4) \times SO(4)$ gauged supergravity was recently investigated in [47], where a large number of holographic renormalization group (RG) flow solutions was also given. In this paper, we will look for supersymmetric Janus solutions in this gauged supergravity.

The study of Janus solutions in $N = 4$ gauged supergravity first appeared in [19], where a number of singular Janus solutions, interpolating between singular geometries, was given. The $N = 4$ gauged supergravity in this case is obtained from a truncation of 11-dimensional supergravity on a tri-Sasakian manifold resulting in a nonsemisimple $SO(3) \ltimes (\mathbf{T}^3, \hat{\mathbf{T}}^3)$ gauge group. In addition, a regular Janus

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solution interpolating between the trivial AdS_4 vacua in $N=4$ gauged supergravity with an $ISO(3)\times ISO(3)$ gauge group, obtained from a nongeometric compactification of type IIB theory, was given in [20]. In this case, the solution involves only scalar fields from the gravity multiplet. Both of these $N=4$ gauged supergravities admit only one supersymmetric AdS_4 vacuum at the origin of the scalar manifold. Therefore, Janus solutions involving more than one critical point are not possible.

Regular Janus solutions, with nonvanishing scalars from both gravity and vector multiplets, in the framework of matter-coupled $N=4$ gauged supergravity with an $SO(4)\times SO(4)$ gauge group appeared only recently in [35]. This $N=4$ gauged supergravity admits a number of supersymmetric AdS_4 vacua [48] and can be obtained from the symplectically deformed $SO(4)\times SO(4)$ gauged supergravity mentioned above for a particular choice of electric-magnetic phases, with two of the phases vanishing and the other two equal to $\frac{\pi}{2}$. However, the solutions found in [35] are obtained only in an $SO(2)\times SO(2)\times SO(3)\times SO(2)$ subtruncation of an $SO(2)\times SO(2)\times SO(2)\times SO(2)$ scalar sector in which only the trivial $SO(4)\times SO(4)$ AdS_4 critical point appears. Accordingly, the Janus solutions in [35] interpolate only between the trivial $SO(4)\times SO(4)$ critical points as well.

In this paper, we will extend this study in two main aspects. We first look at $N=2$ Janus solutions in the full $SO(2)\times SO(2)\times SO(2)\times SO(2)$ scalar sector. Although no free deformation parameters appear in this sector, as shown in [47], we do find a number of new $N=2$ Janus solutions with $SO(2)\times SO(2)\times SO(2)\times SO(2)$ symmetry that generalize the results of [35]. Secondly, we will consider an $SO(3)_{\text{diag}}\times SO(3)$ sector which, in addition to the trivial $SO(4)\times SO(4)$ critical point, admits two nontrivial $N=4$ AdS_4 critical points [47]. We will find $N=1$ supersymmetric Janus solutions that are dependent on the electric-magnetic phases. Moreover, we also find a new family of $N=1$ Janus solutions interpolating between nontrivial AdS_4 critical points. To the best of our knowledge, these are the first Janus solutions that involve nontrivial AdS_4 critical points in the framework of four-dimensional $N=4$ gauged supergravity. Although a large number of Janus solutions of this type can be found in the maximal gauged supergravity (see, for example, [17,21,34,36]), apart from the solutions in three-dimensional $N=8$ gauged supergravity studied recently in [37], no such solutions have been found to date within half-maximal gauged supergravities in higher dimensions. We hope the results of this paper would constitute the first step in filling this gap.

The paper is organized as follows. In Sec. II, we review the structure of four-dimensional $N=4$ gauged supergravity with a symplectically deformed $SO(4)\times SO(4)$ gauge group. We then set up Bogomol'nyi-Prasad-Sommerfield (BPS) equations within $SO(2)\times SO(2)\times SO(2)\times SO(2)$ and $SO(3)_{\text{diag}}\times SO(3)$ truncations, and we find a number of Janus solutions preserving $N=2$ and $N=1$ supersymmetries in Secs. III and IV, respectively. We end the paper with some conclusions and comments in Sec. V.

II. MATTER-COUPLED $N=4$ GAUGED SUPERGRAVITY

In this section, we give a brief review of $N=4$ gauged supergravity coupled to vector multiplets in the embedding tensor formalism constructed in [38]. The gravity and vector multiplets read

$$(e_\mu^{\hat{\mu}}, \psi_\mu^i, A_\mu^m, \chi^i, \tau) \tag{1}$$

and

$$(A_\mu^a, \lambda^{ia}, \phi^{ma}). \tag{2}$$

The bosonic component fields from the gravity and n vector multiplets are given by the graviton $e_\mu^{\hat{\mu}}$, the $6+n$ vector fields $A^{+M}=(A_\mu^m, A_\mu^a)$, a complex scalar τ containing the dilaton ϕ and the axion χ parametrizing the $SL(2,\mathbb{R})/SO(2)$ coset, and $6n$ scalars ϕ^{ma} parametrizing the $SO(6,n)/SO(6)\times SO(n)$ coset. Indices $\mu,\nu,\ldots=0, 1, 2, 3$ and $\hat{\mu},\hat{\nu},\ldots=0, 1, 2, 3$ denote the space-time and tangent space (flat) indices, respectively, while the $m,n=1,\ldots,6$ and $i, j=1, 2, 3, 4$ indices describe fundamental representations of $SO(6)_R$ and $SU(4)_R$ R symmetry. The n vector multiplets are labeled with indices $a,b=1,\ldots,n$. The vector fields A^{+M} and the magnetic dual A^{-M} form a doublet under $SL(2,\mathbb{R})$ and will be collectively denoted by $A^{\alpha M}$, $\alpha=(+,-)$.

The fermionic fields contain four gravitini ψ_μ^i, four spin-$\frac{1}{2}$ fields χ^i, and $4n$ gaugini λ^{ia}. These fields and supersymmetry parameters are subject to the chirality projections

$$\gamma_5\psi_\mu^i=\psi_\mu^i, \qquad \gamma_5\chi^i=-\chi^i, \qquad \gamma_5\lambda^{ia}=\lambda^{ia}, \tag{3}$$

and the same holds for conjugate spinors

$$\gamma_5\psi_{\mu i}=-\psi_{\mu i}, \qquad \gamma_5\chi_i=\chi_i, \qquad \gamma_5\lambda_i^a=-\lambda_i^a. \tag{4}$$

By using the complex scalar τ of the form

$$\tau=\chi+ie^\phi, \tag{5}$$

we can write the coset representative for $SL(2,\mathbb{R})/SO(2)$ as

$$\mathcal{V}_\alpha=e^{\frac{\phi}{2}}\begin{pmatrix}\chi+ie^\phi\\ 1\end{pmatrix}. \tag{6}$$

Similarly, the $6n$ vector multiplet scalars ϕ^{ma} can be described using the coset representative

$$\mathcal{V}_M{}^A = (\mathcal{V}_M{}^m, \mathcal{V}_M{}^a). \tag{7}$$

We have decomposed the $SO(6)\times SO(n)$ index as $A=(m,a)$. We also note that the matrix $\mathcal{V}_M{}^A$ satisfies the relation

$$\eta_{MN} = -\mathcal{V}_M{}^m\mathcal{V}_N{}^m + \mathcal{V}_M{}^a\mathcal{V}_N{}^a, \tag{8}$$

with $\eta_{MN}=\text{diag}(-1,-1,-1,-1,-1,-1,1,\ldots,1)$ being the $SO(6,n)$ invariant tensor. The inverse of $\mathcal{V}_M{}^A$ will be denoted by $\mathcal{V}_A{}^M = (\mathcal{V}_m{}^M, \mathcal{V}_a{}^M)$.

Gaugings of the matter-coupled $N=4$ supergravity are encoded in the components of the embedding tensor $\xi^{\alpha M}$ and $f_{\alpha MNP}$. We will consider only the gaugings with $\xi^{\alpha M}=0$, as required by the existence of supersymmetric AdS_4 vacua [49]. In addition, we will also set all fermionic and vector fields to zero since supersymmetric Janus solutions involve only the metric and scalar fields. The bosonic Lagrangian can then be written as

$$e^{-1}\mathcal{L} = \frac{1}{2}R + \frac{1}{16}\partial_\mu M_{MN}\partial^\mu M^{MN} - \frac{1}{4(\text{Im}\tau)^2}\partial_\mu\tau\partial^\mu\tau^* - V, \tag{9}$$

where $e=\sqrt{-g}$ is the vielbein determinant. The scalar potential is given by

$$V = \frac{1}{16}\left[f_{\alpha MNP}f_{\beta QRS}M^{\alpha\beta}\left[\frac{1}{3}M^{MQ}M^{NR}M^{PS} + \left(\frac{2}{3}\eta^{MQ} - M^{MQ}\right)\eta^{NR}\eta^{PS}\right] - \frac{4}{9}f_{\alpha MNP}f_{\beta QRS}\epsilon^{\alpha\beta}M^{MNPQRS}\right]. \tag{10}$$

The symmetric matrix M_{MN}, with the inverse M^{MN}, is defined by

$$M_{MN} = \mathcal{V}_M{}^m\mathcal{V}_N{}^m + \mathcal{V}_M{}^a\mathcal{V}_N{}^a. \tag{11}$$

The tensor M^{MNPQRS} is obtained from

$$M_{MNPQRS} = \epsilon_{mnpqrs}\mathcal{V}_M{}^m\mathcal{V}_N{}^n\mathcal{V}_P{}^p\mathcal{V}_Q{}^q\mathcal{V}_R{}^r\mathcal{V}_S{}^s \tag{12}$$

by raising the indices with η^{MN}. The matrix $M^{\alpha\beta}$ is the inverse of the symmetric 2×2 matrix $M_{\alpha\beta}$ defined by

$$M_{\alpha\beta} = \text{Re}(\mathcal{V}_\alpha\mathcal{V}_\beta^*). \tag{13}$$

We also need fermionic supersymmetry transformations

$$\delta\psi_\mu^i = 2D_\mu\epsilon^i - \frac{2}{3}A_1^{ij}\gamma_\mu\epsilon_j, \tag{14}$$

$$\delta\chi^i = -\epsilon^{\alpha\beta}\mathcal{V}_\alpha D_\mu\mathcal{V}_\beta\gamma^\mu\epsilon^i - \frac{4}{3}iA_2^{ij}\epsilon_j, \tag{15}$$

$$\delta\lambda_a^i = 2i\mathcal{V}_a{}^M D_\mu\mathcal{V}_M{}^{ij}\gamma^\mu\epsilon_j - 2iA_{2aj}{}^i\epsilon^j, \tag{16}$$

with the fermion shift matrices defined by

$$\begin{aligned} A_1^{ij} &= \epsilon^{\alpha\beta}(\mathcal{V}_\alpha)^*\mathcal{V}_{kl}{}^M\mathcal{V}_N{}^{ik}\mathcal{V}_P{}^{jl}f_{\beta M}{}^{NP}, \\ A_2^{ij} &= \epsilon^{\alpha\beta}\mathcal{V}_\alpha\mathcal{V}_{kl}{}^M\mathcal{V}_N{}^{ik}\mathcal{V}_P{}^{jl}f_{\beta M}{}^{NP}, \\ A_{2ai}{}^j &= \epsilon^{\alpha\beta}\mathcal{V}_\alpha\mathcal{V}_a{}^M\mathcal{V}_{ik}{}^N\mathcal{V}_P{}^{jk}f_{\beta MN}{}^P. \end{aligned} \tag{17}$$

The coset representatives of the forms $\mathcal{V}_M{}^{ij}$ and $\mathcal{V}_{ij}{}^M$ are defined in terms of the 't Hooft symbols G_m^{ij} as

$$\mathcal{V}_M{}^{ij} = \frac{1}{2}\mathcal{V}_M{}^m G_m^{ij} \tag{18}$$

and

$$\mathcal{V}_{ij}{}^M = -\frac{1}{2}\mathcal{V}_m{}^M (G_m^{ij})^*. \tag{19}$$

The explicit representation of G_m^{ij} used in this paper is the same as in [47]. Upper and lower $i,j,\ldots$ indices are related by complex conjugation, as usual.

In this paper, we consider only $N=4$ gauged supergravity coupled to $n=6$ vector multiplets with an $SO(4)\times SO(4)$ gauge group. By decomposing the $SO(6,6)$ fundamental index as $M=(\hat{m},\tilde{m},\hat{a},\tilde{a})$, for $\hat{m},\tilde{m},\hat{a},\tilde{a}=1,2,3$, we can write the embedding tensor for a symplectically deformed $SO(4)\times SO(4)$ gauge group as

$$\begin{aligned} f_{+\hat{m}\hat{n}\hat{p}} &= -g_0\cos\alpha_0\epsilon_{\hat{m}\hat{n}\hat{p}}, & f_{-\hat{m}\hat{n}\hat{p}} &= g_0\sin\alpha_0\epsilon_{\hat{m}\hat{n}\hat{p}}, \\ f_{+\tilde{m}\tilde{n}\tilde{p}} &= -g\cos\alpha\epsilon_{\tilde{m}\tilde{n}\tilde{p}}, & f_{-\tilde{m}\tilde{n}\tilde{p}} &= g\sin\alpha\epsilon_{\tilde{m}\tilde{n}\tilde{p}}, \\ f_{+\hat{a}\hat{b}\hat{c}} &= h_1\cos\beta_1\epsilon_{\hat{a}\hat{b}\hat{c}}, & f_{-\hat{a}\hat{b}\hat{c}} &= h_1\sin\beta_1\epsilon_{\hat{a}\hat{b}\hat{c}}, \\ f_{+\tilde{a}\tilde{b}\tilde{c}} &= h_2\cos\beta_2\epsilon_{\tilde{a}\tilde{b}\tilde{c}}, & f_{-\tilde{a}\tilde{b}\tilde{c}} &= h_2\sin\beta_2\epsilon_{\tilde{a}\tilde{b}\tilde{c}}. \end{aligned} \tag{20}$$

These components of the embedding tensor were given in [50], and we have rewritten them in the notation of [42]. $f_{\pm\hat{m}\hat{n}\hat{p}}$ and $f_{\pm\tilde{m}\tilde{n}\tilde{p}}$ describe the embedding of the first $SO(4)\sim SO(3)\times SO(3)$ factor in $SO(6)_R$ R symmetry. As previously mentioned, the constants α_0 and α can be set to zero and $\frac{\pi}{2}$, respectively. g_0, g, h_1, and h_2 are gauge coupling constants for the four $SO(3)$ factors. In subsequent sections, we will look for supersymmetric Janus solutions with different numbers of unbroken supersymmetries and residual symmetries.

III. $N=2$ SUPERSYMMETRIC JANUS SOLUTIONS

We begin with a truncation to scalars that are singlets of the $SO(2)\times SO(2)\times SO(2)\times SO(2)$ subgroup of the $SO(4)\times SO(4)$ gauge group. We first choose an explicit form of $SO(6,6)$ generators in the fundamental representation as

$$(t_{MN})_P{}^Q = 2\delta^Q_{[M}\eta_{N]P}. \quad (21)$$

The $SO(6,6)$ noncompact generators are accordingly given by

$$Y_{ma} = t_{m,a+6}. \quad (22)$$

Following [35], one can write the coset representative for $SO(2)\times SO(2)\times SO(2)\times SO(2)$ singlet scalars as

$$\mathcal{V} = e^{\phi_1 Y_{33}} e^{\phi_2 Y_{36}} e^{\phi_3 Y_{63}} e^{\phi_4 Y_{66}}. \quad (23)$$

The metric ansatz takes the form of the usual AdS$_3$-sliced domain walls

$$ds^2 = e^{2A(r)}\left(e^{\frac{2\rho}{\ell}} dx^2_{1,1} + d\rho^2\right) + dr^2, \quad (24)$$

in which ℓ denotes the radius of the AdS$_3$ slices. $dx^2_{1,1} = \eta_{\alpha\beta}dx^\alpha dx^\beta$, α, $\beta = 0$, 1, is the flat metric on two-dimensional Minkowski space.

All scalars ϕ_i, $i = 1, 2, 3, 4$, together with the dilaton ϕ and the axion χ, are allowed to depend only on r. Analyses of relevant BPS equations have already appeared in many places (see, for example, [17,18]), so we will simply summarize the results. The supersymmetry transformations $\delta\psi^i_{\hat{\alpha}}$ give the following equation,

$$A'^2 = W^2 - \frac{1}{\ell^2} e^{-2A}, \quad (25)$$

while $\delta\psi^i_{\hat{\rho}}$ gives the Killing spinor of the form

$$\epsilon^{\hat{i}} = e^{\frac{\rho}{2\ell}} \tilde{\epsilon}^{\hat{i}} \quad (26)$$

for ρ-independent spinors $\tilde{\epsilon}^{\hat{i}}$. In Eq. (25), $W = |\mathcal{W}|$, and the superpotential

$$\mathcal{W} = \frac{2}{3}\hat{\alpha} \quad (27)$$

is obtained from the eigenvalue $\hat{\alpha}$ of A_1^{ij} with the corresponding eigenvectors $\epsilon^{\hat{i}}$ identified with the Killing spinors. We use an index $\hat{i}$ to count the number of unbroken supersymmetries.

With the projectors

$$\gamma_{\hat{r}}\epsilon^{\hat{i}} = e^{i\Lambda}\epsilon_{\hat{i}} \quad (28)$$

and

$$\gamma_{\hat{\rho}}\epsilon^{\hat{i}} = i\kappa e^{i\Lambda}\epsilon_{\hat{i}}, \quad (29)$$

with $\kappa^2 = 1$ and an r-dependent phase Λ, the Killing spinors can be determined from $\delta\psi^i_{\hat{r}}$ to be

$$\epsilon^{\hat{i}} = e^{\frac{A}{2}+\frac{\rho}{2\ell}+i\frac{\Lambda}{2}}\varepsilon^{(0)\hat{i}}. \quad (30)$$

The spinors $\varepsilon^{(0)\hat{i}}$ can (possibly) have an r-dependent phase and satisfy the following projection conditions:

$$\gamma_{\hat{r}}\varepsilon^{(0)\hat{i}} = \varepsilon^{(0)}_{\hat{i}}, \qquad \gamma_{\hat{\rho}}\varepsilon^{(0)\hat{i}} = i\kappa\varepsilon^{(0)}_{\hat{i}}. \quad (31)$$

With all these results, the conditions $\delta\psi^i_{\hat{\alpha}}$ determine the explicit form of the phase $e^{i\Lambda}$ to be

$$e^{i\Lambda} = \frac{\mathcal{W}}{A' + \frac{i\kappa}{\ell}e^{-A}} = \frac{\mathcal{W}}{W^2}\left(A' - \frac{i\kappa}{\ell}e^{-A}\right). \quad (32)$$

With the projector (28), the variations $\delta\chi^i$ and $\delta\lambda^i_a$ lead to the BPS equations for scalars. Finally, we note that the sign factor $\kappa = \pm 1$ corresponds to chiralities of the Killing spinors on the two-dimensional defects.

For the $SO(2)\times SO(2)\times SO(2)\times SO(2)$ truncation, the A_1^{ij} tensor takes the form (see [47] for more details)

$$A_1^{ij} = \mathrm{diag}(\mathcal{A}_-, \mathcal{A}_+, \mathcal{A}_+, \mathcal{A}_-). \quad (33)$$

Both of the eigenvalues lead to an $N = 2$ unbroken supersymmetry with the superpotential $\mathcal{W}_\mp = \frac{2}{3}\mathcal{A}_\mp$ and the Killing spinors $\epsilon^{1,4}$ and $\epsilon^{2,3}$, respectively. Following [47], we will set $\epsilon^2 = \epsilon^3 = 0$ and choose the superpotential to be

$$\begin{aligned}\mathcal{W} &= \mathcal{W}_- \\ &= \frac{1}{2}e^{-\frac{\phi}{2}}[\cosh\phi_4[g\cosh\phi_3(e^\phi\sin\alpha + i\cos\alpha) - g_0\sinh\phi_1\sinh\phi_3] \\ &\quad -g_0\cosh\phi_1(\cosh\phi_2 + i\sinh\phi_2\sinh\phi_4) + ig\sin\alpha\cosh\phi_3\cosh\phi_4\chi]. \end{aligned} \quad (34)$$

The scalar potential can be written in terms of the superpotential as

$$\begin{aligned} V &= -2G^{rs}\frac{\partial W}{\partial \Phi^r}\frac{\partial W}{\partial \Phi^s} - 3W^2 \\ &= -\frac{1}{4}e^{-\phi}[g^2(1+\cos 2\alpha) + 2g_0^2 + 2g^2 \sin\alpha\chi(2\cos\alpha + \sin\alpha\chi)] - \frac{1}{2}e^{\phi}g^2\sin^2\alpha \\ &\quad + 2gg_0 \sin\alpha \cosh\phi_1 \cosh\phi_2 \cosh\phi_3 \cosh\phi_4, \end{aligned} \tag{35}$$

in which we have defined the scalars $\Phi^r = (\phi, \chi, \phi_1, \phi_2, \phi_3, \phi_4)$. G^{rs} is the inverse of the scalar metric appearing in the scalar kinetic terms.

With the coset representative (23), the kinetic term for scalar fields is given by

$$\begin{aligned} \mathcal{L}_{\text{kin}} &= \frac{1}{2}G_{rs}\Phi^{r\prime}\Phi^{s\prime} \\ &= -\frac{1}{4}(\phi'^2 + e^{-2\phi}\chi'^2) - \frac{1}{16}[6 + \cosh 2(\phi_2 - \phi_3) \\ &\quad + \cosh 2(\phi_2 + \phi_3) + 2\cosh 2\phi_4(\cosh 2\phi_2 \cosh 2\phi_3 - 1)]\phi_1'^2 \\ &\quad - \cosh\phi_2 \cosh\phi_4 \sinh\phi_3 \sinh\phi_4 \phi_1'\phi_2' - \cosh\phi_3 \cosh\phi_4 \sinh\phi_2 \sinh\phi_4 \phi_1'\phi_3' \\ &\quad + \sinh\phi_2 \sinh\phi_3 \phi_1'\phi_4' - \frac{1}{2}\cosh^2\phi_4 \phi_2'^2 - \frac{1}{2}\cosh^2\phi_4 \phi_3'^2 - \frac{1}{2}\phi_4'^2, \end{aligned} \tag{36}$$

from which we can determine the scalar metric G_{rs} and its inverse G^{rs}. Since G^{rs} will appear in the final form of the BPS equations, for later convenience we will give its explicit form here,

$$G^{rs} = \begin{pmatrix} -2 & 0 & \mathbf{0}_{1\times 4} \\ 0 & -2e^{2\phi} & \mathbf{0}_{1\times 4} \\ \mathbf{0}_{4\times 1} & \mathbf{0}_{4\times 1} & \hat{G}^{\hat{r}\hat{s}} \end{pmatrix}, \tag{37}$$

with the 4×4 symmetric matrix $\hat{G}^{\hat{r}\hat{s}}$, for $\hat{r}, \hat{s} = 1, 2, 3, 4$, given by

$$\hat{G}^{\hat{r}\hat{s}} = \begin{pmatrix} \Box_1 & \Delta_1 & \Delta_2 & \Delta_3 \\ \Delta_1 & \Box_2 & \Delta_4 & \Delta_5 \\ \Delta_2 & \Delta_4 & \Box_3 & \Delta_6 \\ \Delta_3 & \Delta_5 & \Delta_6 & \Box_4 \end{pmatrix} \tag{38}$$

and

$$\begin{aligned} &\Box_1 = -\text{sech}^2\phi_2 \text{sech}^2\phi_3, \qquad \Box_2 = -\text{sech}^2\phi_3 \text{sech}^2\phi_4 - \tanh^2\phi_3, \\ &\Box_3 = \text{sech}^2\phi_2 \tanh^2\phi_4 - 1, \qquad \Box_4 = -\frac{1}{2}\text{sech}^2\phi_2 \text{sech}^2\phi_3(1 + \cosh 2\phi_2 \cosh 2\phi_3), \\ &\Delta_1 = \text{sech}\phi_2 \text{sech}\phi_3 \tanh\phi_3 \tanh\phi_4, \qquad \Delta_2 = \text{sech}\phi_2 \text{sech}\phi_3 \tanh\phi_2 \tanh\phi_4, \\ &\Delta_3 = -\text{sech}\phi_2 \text{sech}\phi_3 \tanh\phi_2 \tanh\phi_3, \qquad \Delta_4 = -\tanh\phi_2 \tanh\phi_3 \tanh^2\phi_4, \\ &\Delta_5 = \tanh\phi_2 \tanh^2\phi_3 \tanh\phi_4, \qquad \Delta_6 = \tanh^2\phi_2 \tanh\phi_3 \tanh\phi_4. \end{aligned} \tag{39}$$

The scalar potential and superpotential admit one AdS_4 critical point at $\phi_1 = \phi_2 = \phi_3 = \phi_4 = 0$ and

$$\phi = \ln\left[-\frac{g_0}{g\sin\alpha}\right], \qquad \chi = -\frac{\cos\alpha}{\sin\alpha}. \tag{40}$$

By shifting the dilaton and axion, or equivalently by choosing $g_0 = -g$ for $\alpha = \frac{\pi}{2}$, we can bring this critical point to the origin of the scalar manifold $SL(2,\mathbb{R})/SO(2) \times SO(6,6)/SO(6) \times SO(6)$, at which all scalars vanish. With this choice, the cosmological constant and AdS_4 radius are given by

$$V_0 = -3g^2, \qquad L = \sqrt{-\frac{3}{V_0}} = \frac{1}{g}, \tag{41}$$

in which we have taken $g > 0$ without loss of generality. This critical point is invariant under the full $SO(4)\times SO(4)$ gauge symmetry and preserves $N = 4$ supersymmetry.

Using the projector (28) and the superpotential (34), we find that all the BPS conditions with $\epsilon^{2,3} = 0$ lead to the following BPS equations:

$$A'^2 + \frac{1}{\ell^2}e^{-2A} = W^2, \tag{42}$$

$$\phi' = -4\frac{A'}{W}\frac{\partial W}{\partial\phi} - 4e^{\phi}\frac{\kappa e^{-A}}{\ell W}\frac{\partial W}{\partial\chi}, \tag{43}$$

$$\chi' = -4e^{2\phi}\frac{A'}{W}\frac{\partial W}{\partial\chi} + 4e^{\phi}\frac{\kappa e^{-A}}{\ell W}\frac{\partial W}{\partial\phi}, \tag{44}$$

$$\phi_1' = \hat{G}^{1\hat{r}}\frac{A'}{W}\frac{\partial W}{\partial\hat{\Phi}^{\hat{r}}} - 2\mathrm{sech}\phi_2\mathrm{sech}\phi_3\mathrm{sech}\phi_4\frac{\kappa e^{-A}}{\ell W}\frac{\partial W}{\partial\phi_3}, \tag{45}$$

$$\phi_2' = \hat{G}^{2\hat{r}}\frac{A'}{W}\frac{\partial W}{\partial\hat{\Phi}^{\hat{r}}} + \frac{\kappa e^{-A}}{\ell W}\left(2\mathrm{sech}\phi_4\tanh\phi_3\tanh\phi_4\frac{\partial W}{\partial\phi_3} - 2\mathrm{sech}\phi_4\frac{\partial W}{\partial\phi_4}\right), \tag{46}$$

$$\phi_3' = \hat{G}^{3\hat{r}}\frac{A'}{W}\frac{\partial W}{\partial\hat{\Phi}^{\hat{r}}} + \frac{\kappa e^{-A}}{\ell W}\left(2\mathrm{sech}\phi_2\mathrm{sech}\phi_3\mathrm{sech}\phi_4\frac{\partial W}{\partial\phi_1}\right.$$
$$\left.-2\mathrm{sech}\phi_4\tanh\phi_3\tanh\phi_4\frac{\partial W}{\partial\phi_2} + 2\mathrm{sech}\phi_4\tanh\phi_2\tanh\phi_3\frac{\partial W}{\partial\phi_4}\right), \tag{47}$$

$$\phi_4' = \hat{G}^{4\hat{r}}\frac{A'}{W}\frac{\partial W}{\partial\hat{\Phi}^{\hat{r}}} + \frac{\kappa e^{-A}}{\ell W}\left(2\mathrm{sech}\phi_4\frac{\partial W}{\partial\phi_2} - 2\mathrm{sech}\phi_4\tanh\phi_2\tanh\phi_3\frac{\partial W}{\partial\phi_3}\right), \tag{48}$$

with $\hat{\Phi}^{\hat{r}} = (\phi_1, \phi_2, \phi_3, \phi_4)$. Before giving the solutions, we first note that in the limit $\ell \to \infty$, these equations reduce to the BPS equations for RG flows studied in [47], as expected. Furthermore, for $\phi_2 = \phi_4 = 0$ or $\phi_1 = \phi_3 = 0$, we recover the BPS equations for the Janus solutions with $SO(2)\times SO(2)\times SO(2)\times SO(3)$ or $SO(2)\times SO(2)\times SO(3)\times SO(2)$ symmetries studied in [35].

We now give $N = 2$ supersymmetric Janus solutions with $SO(2)\times SO(2)\times SO(2)\times SO(2)$ symmetry. After numerically solving the BPS equations, we find examples of Janus solutions for $g=1$, $\kappa=1$, $\ell=1$ and $g_0 = -g\sin\alpha$, as in Fig. 1. In the figure, we have depicted the solutions for different values of the phase α. We also emphasize here that all values of α are equivalent to $\alpha = \frac{\pi}{2}$. We have given the solutions for various values of α only for clarity of presentation since solutions with different boundary conditions but the same value of α are very close to each other and difficult to see. These solutions interpolate between $SO(4)\times SO(4)$ critical points and describe two-dimensional conformal defects within the $N = 4$ SCFT. The defects are invariant under the $SO(2)\times SO(2)\times SO(2)\times SO(2)$ subgroup of the $SO(4)\times SO(4)$ symmetry of the three-dimensional SCFT and preserve $N = (2,0)$ or $N = (0,2)$ supersymmetry in two dimensions, depending on whether the value of κ is 1 or -1.

IV. $N = 1$ SUPERSYMMETRIC JANUS SOLUTIONS

We now move to the $SO(3)_{\mathrm{diag}}\times SO(3)$ sector, which is a subtruncation of the $SO(3)_{\mathrm{diag}}$ sector studied in [47]. We will follow the notation of [47] for the sake of comparison. The $SO(3)_{\mathrm{diag}}\times SO(3)$ sector contains two singlet scalars from an $SO(6,6)/SO(6)\times SO(6)$ coset (see [47] for additional details), with the coset representative

$$\mathcal{V} = e^{\phi_1\hat{Y}_1}e^{\phi_3\hat{Y}_3}, \tag{49}$$

in which the noncompact generators are given by

$$\hat{Y}_1 = Y_{11} + Y_{22} + Y_{33} + Y_{44},$$
$$\hat{Y}_3 = Y_{51} + Y_{62} + Y_{73} + Y_{84}. \tag{50}$$

The A_1^{ij} tensor takes the form

$$A_1^{ij} = \mathrm{diag}(\mathcal{A}, \mathcal{B}, \mathcal{B}, \mathcal{B}), \tag{51}$$

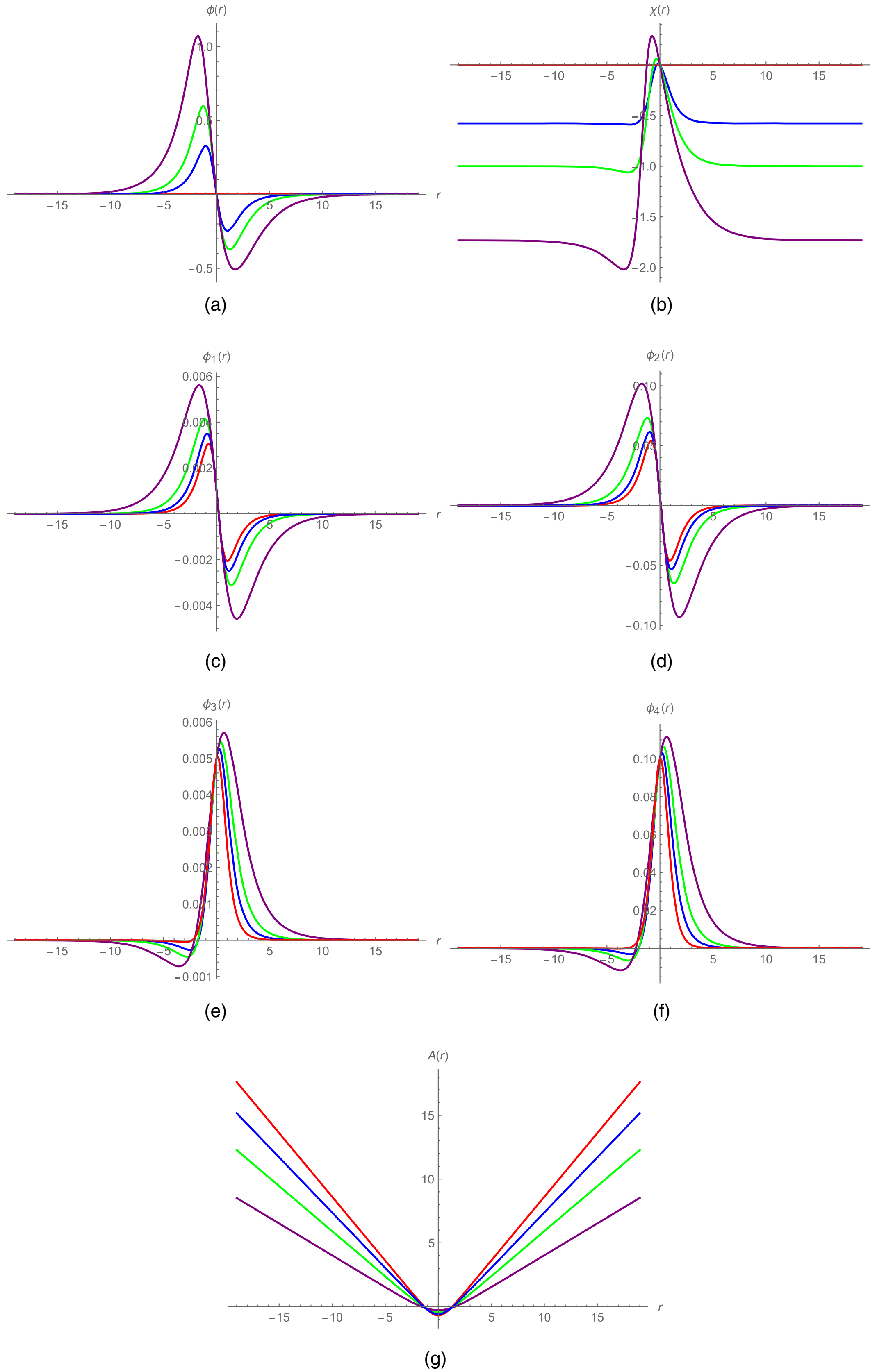

FIG. 1. Examples of $N = 2$ Janus solutions interpolating between $N = 4$ AdS_4 critical points with $SO(4) \times SO(4)$ symmetry for $g = 1$, $\kappa = 1$, $\ell = 1$, $g_0 = -g \sin\alpha$ and $\alpha = \frac{\pi}{6}$ (purple lines), $\alpha = \frac{\pi}{4}$ (green lines), $\alpha = \frac{\pi}{3}$ (blue lines), $\alpha = \frac{\pi}{2}$ (red lines). (a) $\phi(r)$ solution. (b) $\chi(r)$ solution. (c) $\phi_1(r)$ solution. (d) $\phi_2(r)$ solution. (e) $\phi_3(r)$ solution. (f) $\phi_4(r)$ solution. (g) $A(r)$ solution.

with $\mathcal{A}$ leading to the superpotential

$$\mathcal{W} = \frac{1}{2}e^{\frac{\phi}{2}}[g\cosh^3\phi_3 + h_1\sin\beta_1(i\sinh\phi_1 - \cosh\phi_1\sinh\phi_3)^3] + \frac{1}{2}e^{-\frac{\phi}{2}}[g(\cosh\phi_1 + i\sinh\phi_1\sinh\phi_3)^3 - (\sinh\phi_1 + i\cosh\phi_1\sinh\phi_3)^3 \times h_1\cos\beta_1] + \frac{1}{2}e^{-\frac{\phi}{2}}[ig\cosh^3\phi_3 + h_1\sin\beta_1(\sinh\phi_1 + i\cosh\phi_1\sinh\phi_3)^3]\chi. \quad (52)$$

The solutions in this sector then preserve $N=1$ supersymmetry. To simplify the expressions, in this case we will set $\alpha = \frac{\pi}{2}$ and $g_0 = -g$.

For completeness, we also note that the scalar potential can be written as

$$V = 4\left(\frac{\partial W}{\partial\phi}\right)^2 + 4e^{2\phi}\left(\frac{\partial W}{\partial\chi}\right)^2 + \frac{2}{3}\mathrm{sech}^2\phi_3\left(\frac{\partial W}{\partial\phi_1}\right)^2 + \frac{2}{3}\left(\frac{\partial W}{\partial\phi_3}\right)^2 - 3W^2. \quad (53)$$

The explicit form of this potential can be found in [47]. In this paper, we simply recall that the scalar potential admits three supersymmetric AdS_4 critical points. The first one is the trivial $SO(4)\times SO(4)$ critical point at which all scalars vanish for $\alpha = \frac{\pi}{2}$ and $g_0 = -g$, while the other two are given by

$$i:\ \beta_1 = 0; \quad \phi_3 = \chi = 0, \quad \phi_1 = \frac{1}{2}\ln\left[\frac{h_1+g}{h_1-g}\right], \quad \phi = -\frac{1}{2}\ln\left[1-\frac{g^2}{h_1^2}\right], \quad V_0 = -\frac{3g^2h_1}{\sqrt{h_1^2-g^2}}, \quad (54)$$

$$ii:\ \beta_1 = \frac{\pi}{2}; \quad \phi_1 = \chi = 0, \quad \phi_3 = \frac{1}{2}\ln\left[\frac{h_1+g}{h_1-g}\right], \quad \phi = \frac{1}{2}\ln\left[1-\frac{g^2}{h_1^2}\right], \quad V_0 = -\frac{3g^2h_1}{\sqrt{h_1^2-g^2}}. \quad (55)$$

Both of these critical points preserve $N=4$ supersymmetry, as can be verified by setting $\chi = \phi_1 = 0$ or $\chi = \phi_3 = 0$, which gives $\mathcal{A} = \mathcal{B}$. On the other hand, for $\phi_1 \neq 0$ and $\phi_3 \neq 0$, the supersymmetry is broken to $N=1$. The holographic RG flows between these critical points preserving $N=4$ and $N=1$ supersymmetries were already studied in [47].

In this work, we are interested in supersymmetric Janus solutions. We first note that setting either $\phi_1 = 0$ or $\phi_3 = 0$ does not lead to a consistent set of BPS equations for Janus solutions. This implies that, unlike in the RG flow case, there are no $N=4$ supersymmetric Janus solutions with $SO(3)_{\mathrm{diag}}\times SO(3)\times SO(3)$ or $SO(3)\times SO(3)_{\mathrm{diag}}\times SO(3)$ symmetries. For $\phi_1 \neq 0$ and $\phi_3 \neq 0$, truncating out χ is also not consistent with the BPS equations. Therefore, $N=1$ Janus solutions must involve all scalars in the $SO(3)_{\mathrm{diag}}\times SO(3)$ sector, as in the case of the $N=1$ RG flow solutions found in [47].

Using the same procedure as in the previous section with $\epsilon^2 = \epsilon^3 = \epsilon^4 = 0$, we find that the BPS equations can be written as

$$\phi' = -4\frac{A'}{W}\frac{\partial W}{\partial\phi} - 4e^{\phi}\frac{e^{-A}\kappa}{\ell W}\frac{\partial W}{\partial\chi}, \quad (56)$$

$$\chi' = -4e^{2\phi}\frac{A'}{W}\frac{\partial W}{\partial\phi} + 4e^{\phi}\frac{e^{-A}\kappa}{\ell W}\frac{\partial W}{\partial\phi}, \quad (57)$$

$$\phi_1' = -\frac{2}{3}\mathrm{sech}^2\phi_3\frac{A'}{W}\frac{\partial W}{\partial\phi_1} - \frac{2}{3}\mathrm{sech}\phi_3\frac{e^{-A}\kappa}{\ell W}\frac{\partial W}{\partial\phi_3}, \quad (58)$$

$$\phi_3' = -\frac{2}{3}\phi_3\frac{A'}{W}\frac{\partial W}{\partial\phi_3} + \frac{2}{3}\mathrm{sech}\phi_3\frac{e^{-A}\kappa}{\ell W}\frac{\partial W}{\partial\phi_1} \quad (59)$$

together with the usual equation for the metric function

$$A'^2 + \frac{e^{-2A}}{\ell^2} = W^2, \quad (60)$$

with the superpotential given in (52). It should be noted that these equations again reduce to the BPS equations for holographic RG flows studied in [47] in the limit $\ell \to \infty$, as expected.

We begin with generic solutions for different values of the phase β_1. After numerically solving the BPS equations, we find examples of solutions for $g=1$, $h_1=2$, $\kappa=1$, and $\ell=1$, as shown in Fig. 2. Some of the solutions are very close to each other, so some solutions, in particular the one represented by the green lines, are not clearly seen. In the figure, all the solutions are qualitatively similar and describe two-dimensional conformal defects within a three-dimensional $N=4$ SCFT with $SO(4)\times SO(4)$ symmetry. Unlike the solutions in the previous section, these defects preserve only $N=(1,0)$ or $N=(0,1)$ supersymmetry, depending on whether the value of κ is 1 or -1.

For $\beta_1 = 0$, there are two AdS_4 critical points, the trivial one and critical point (i). With appropriate boundary

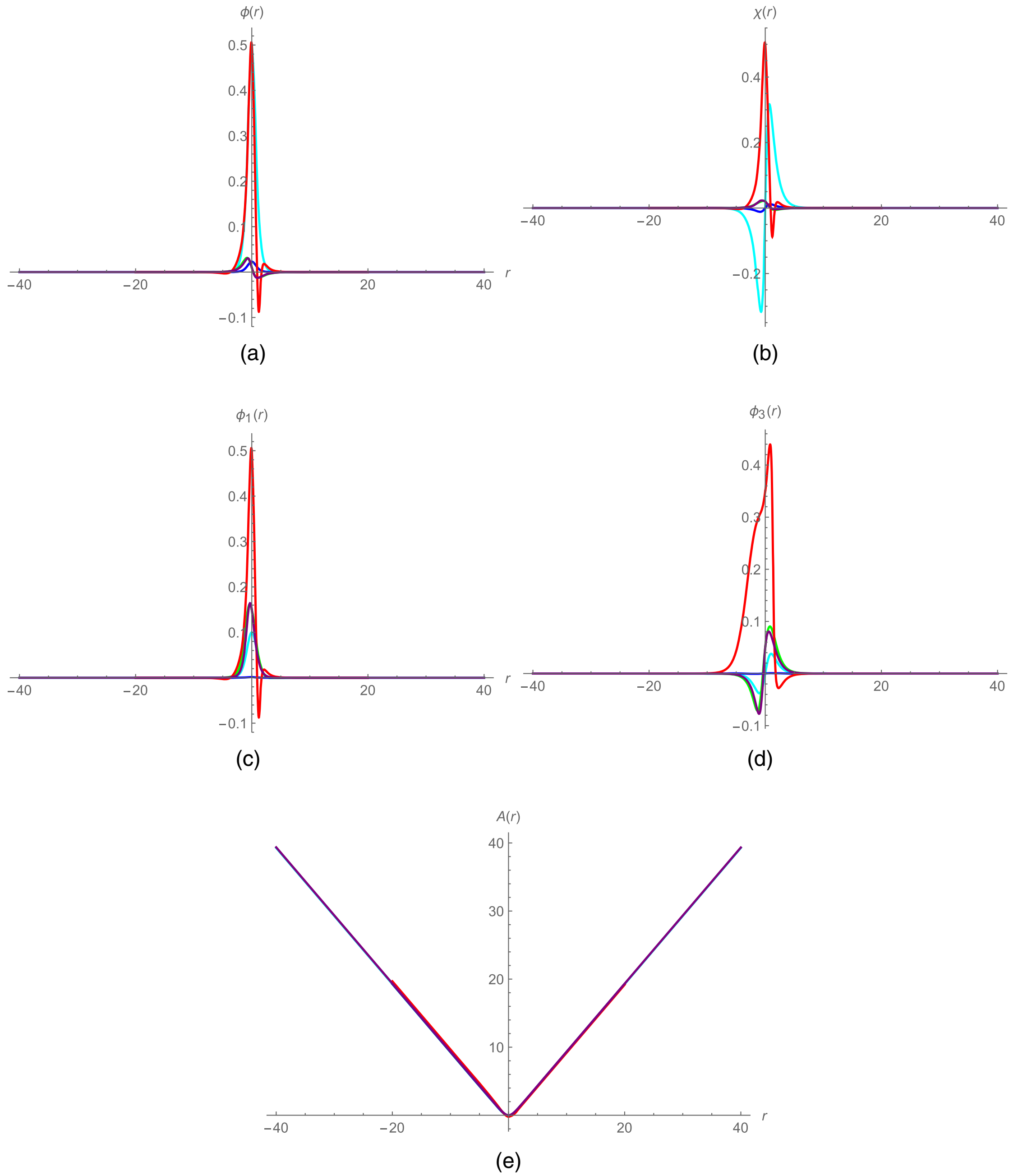

FIG. 2. Examples of $N = 1$ Janus solutions interpolating between $N = 4$ AdS_4 critical points with $SO(4) \times SO(4)$ symmetry for different values of β_1, $\beta_1 = 0$ (cyan lines), $\beta_1 = \frac{\pi}{6}$ (red lines), $\beta_1 = \frac{\pi}{4}$ (green lines), $\beta_1 = \frac{\pi}{3}$ (blue lines), $\beta_1 = \frac{\pi}{2}$ (purple lines). (a) $\phi(r)$ solution. (b) $\chi(r)$ solution. (c) $\phi_1(r)$ solution. (d) $\phi_3(r)$ solution. (e) $A(r)$ solution.

conditions, we find a Janus solution interpolating between critical point (i) for $g = 1$, $h_1 = 2$, $\kappa = 1$, and $\ell = 1$, as shown in Fig. 3. This solution is represented by the pink lines. We have included the Janus solution between $SO(4) \times SO(4)$ critical points (cyan lines) for comparison. We have also given the solution for $A'(r)$ to explicitly show that the two solutions indeed interpolate between different pairs of critical points. However, it should be noted that the critical point (i) on both sides is generated by a holographic RG flow from the $SO(4) \times SO(4)$ critical point. In particular, this RG flow was one of the solutions studied recently in [47]. The Janus solution is accordingly similar to those given in [17,21,34,36,37]. A similar solution interpolating between critical points (ii) can also be found, as shown by the yellow lines in Fig. 4 with $g = 1$, $h_1 = 2$, $\kappa = 1$, and $\ell = 1$. As in Fig. 3, we have included the Janus solution between the $SO(4) \times SO(4)$ critical points for comparison (purple lines). These two solutions describe $N = (1, 0)$ or $N = (0, 1)$ conformal defects within $N = 4$ SCFTs with $SO(3)_{\mathrm{diag}} \times SO(3) \times SO(3)$ or $SO(3) \times SO(3)_{\mathrm{diag}} \times SO(3)$ symmetry.

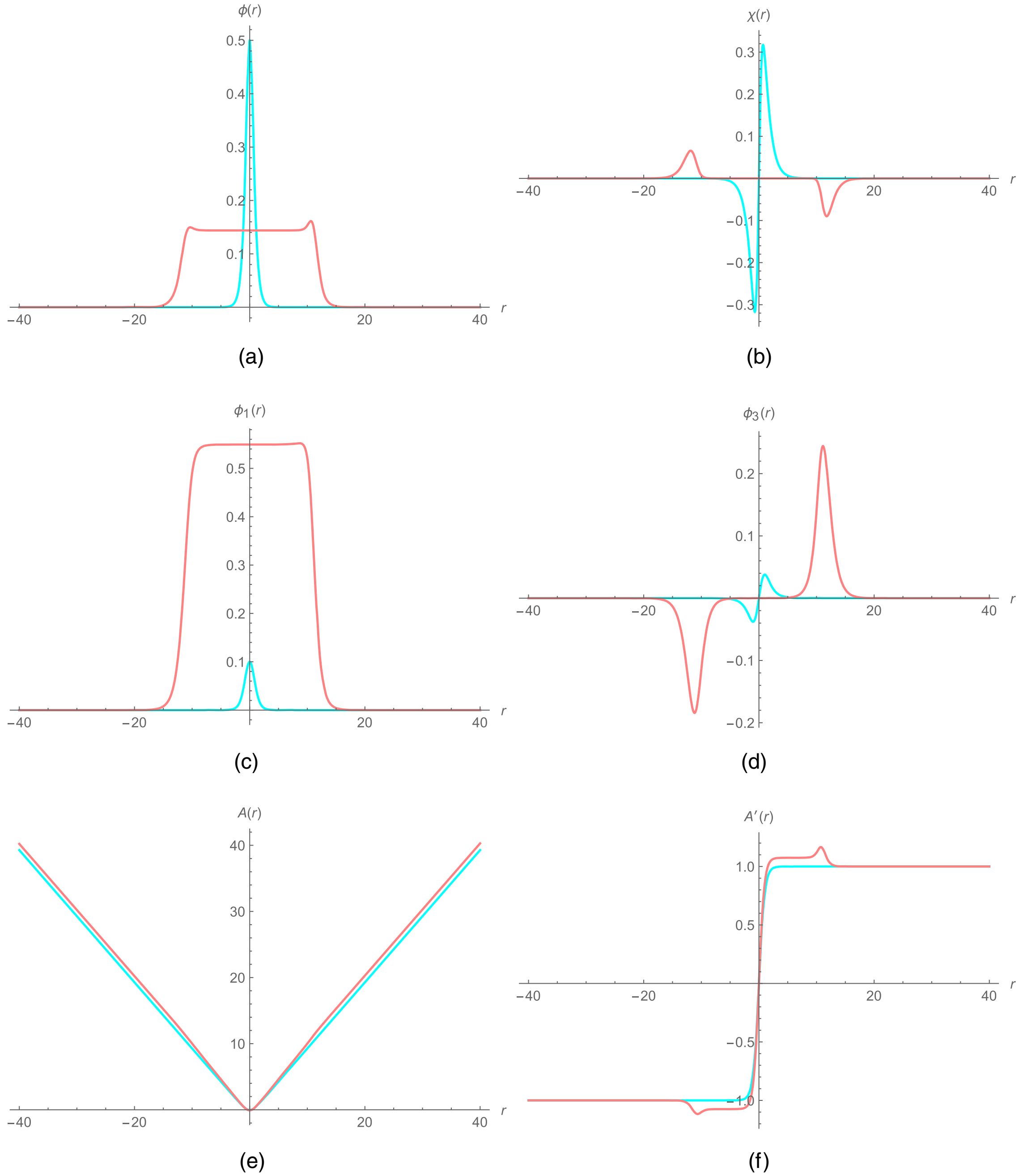

FIG. 3. An example of $N = 1$ Janus solutions (pink lines) interpolating between $N = 4$ AdS_4 critical points with $SO(3)_{\rm diag} \times SO(3) \times SO(3)$ symmetry [critical point (i)]. (a) $\phi(r)$ solution. (b) $\chi(r)$ solution. (c) $\phi_1(r)$ solution. (d) $\phi_3(r)$ solution. (e) $A(r)$ solution. (f) $A'(r)$ solution.

V. CONCLUSIONS AND DISCUSSIONS

In this paper, we have studied supersymmetric Janus solutions using four-dimensional $N = 4$ gauged supergravity with an $SO(4) \times SO(4)$ gauge group in the presence of symplectic deformations. We have found two classes of solutions preserving $N = 1$ and $N = 2$ supersymmetries. The $N = 2$ solutions interpolate between the trivial $N = 4$ critical points with $SO(4) \times SO(4)$ symmetry. In this case, electric-magnetic phases or deformation parameters do not appear, apart from those fixed by $SL(2, \mathbb{R})$ transformations and redefinitions of the dilaton and axion, and there are no other AdS_4 critical points. The solutions are invariant under $SO(2) \times SO(2) \times SO(2) \times SO(2)$ symmetry and describe $N = (2, 0)$ or $N = (0, 2)$ two-dimensional conformal defects in the $N = 4$ SCFT dual to the AdS_4 critical point.

On the other hand, in the $N = 1$ case, we have found more interesting solutions. The solutions are obtained in the $SO(3)_{\rm diag} \times SO(3)$ sector, and for the particular values of

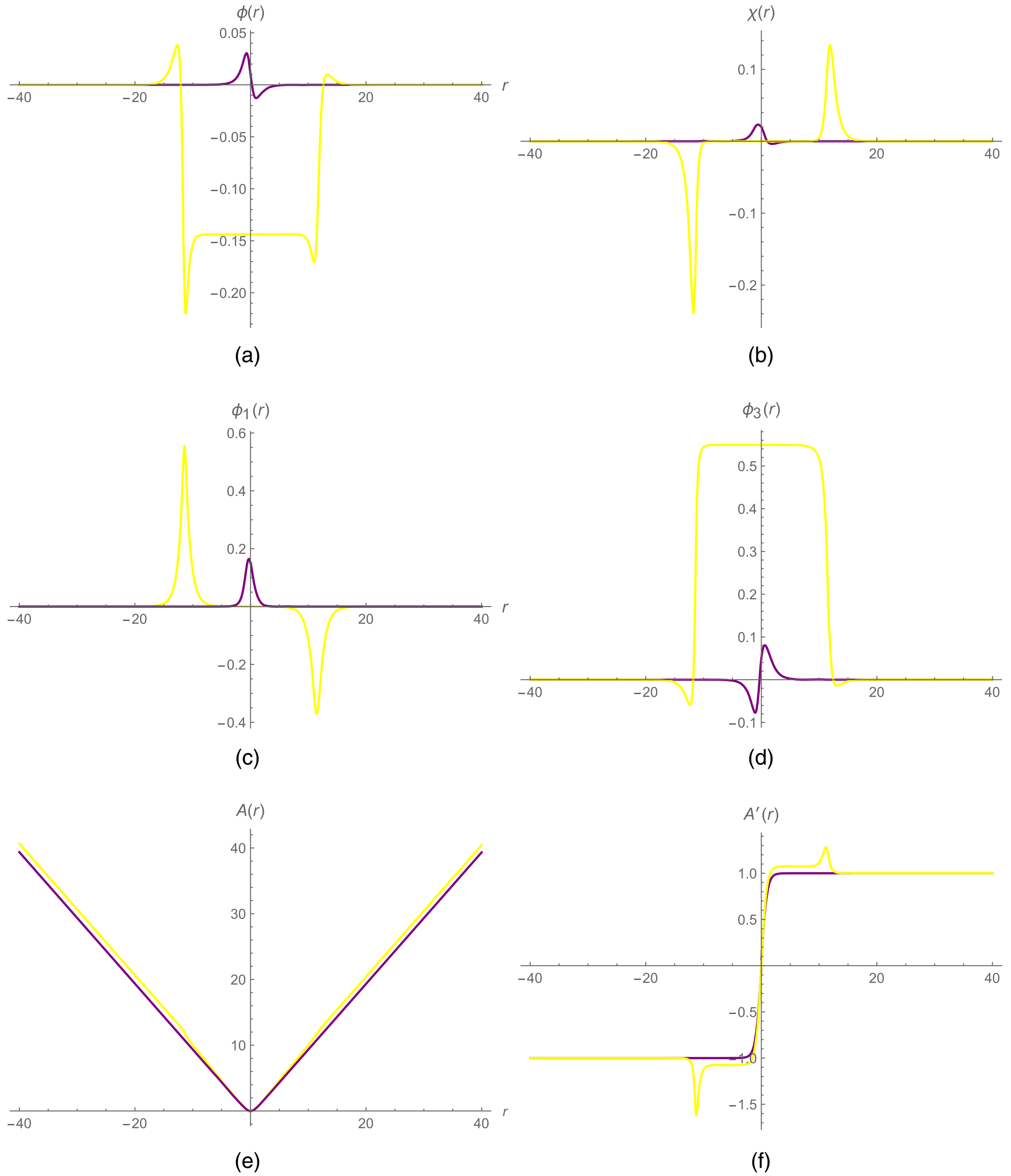

FIG. 4. An example of $N = 1$ Janus solutions (yellow lines) interpolating between $N = 4$ AdS_4 critical points with $SO(3) \times SO(3)_{\mathrm{diag}} \times SO(3)$ symmetry [critical point (ii)]. (a) $\phi(r)$ solution. (b) $\chi(r)$ solution. (c) $\phi_1(r)$ solution. (d) $\phi_3(r)$ solution. (e) $A(r)$ solution. (f) $A'(r)$ solution.

the phase $\beta_1 = 0$ and $\beta_1 = \frac{\pi}{2}$, there are two additional nontrivial $N = 4$ critical points with $SO(3)_{\mathrm{diag}} \times SO(3) \times SO(3)$ and $SO(3) \times SO(3)_{\mathrm{diag}} \times SO(3)$ symmetries apart from the trivial critical point. There are $N = 1$ solutions interpolating between $SO(4) \times SO(4)$ critical points for any values of the electric-magnetic phase β_1, as in the $N = 2$ solutions. Moreover, we have found solutions interpolating between $SO(3)_{\mathrm{diag}} \times SO(3) \times SO(3)$ critical points and between $SO(3) \times SO(3)_{\mathrm{diag}} \times SO(3)$ critical points. In this case, the solutions describe two-dimensional conformal defects in $N = 4$ SCFTs dual to AdS_4 critical points (i) and (ii) that preserve $N = (1, 0)$ or $N = (0, 1)$ supersymmetries on the defects. These are the first examples of Janus solutions in $N = 4$ gauged supergavity that involve nontrivial AdS_4 critical points.

It would be interesting to identify the $N = 4$ SCFTs dual to the AdS_4 critical points considered here and study the conformal defects dual to the Janus solutions found in this paper. As pointed out in [47], in the $SO(3)_{\rm diag}$ invariant scalar sector, both of the electric-magnetic phases β_1 and β_2 appear in the scalar potential and the superpotential. It would be of particular interest to investigate this sector and look for new supersymmetric AdS_4 vacua and to also find new Janus solutions in this case. Finally, since $SO(4)\times SO(4)$ gauged supergravity admitting AdS_4 vacua for any values of the deformation parameters presently has no known embedding in higher dimensions, it would be highly desirable to find the corresponding embedding that would provide an uplift for the solutions found here and those given in [35,47,48] to 10/11 dimensions. Along these lines, recent developments in the double field theory formalism would be very useful; see, for example, [50–56]. The uplifted solutions should provide a complete gravity dual of the $N=4$ SCFTs in three dimensions together with deformations and conformal defects in a string/M-theory context. We leave these issues for future work.

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