PHYSICAL REVIEW D **85,** 035003 (2012)

Possibility of generating leading-order gaugino masses in a direct gauge mediation scenario

Tirtha Sankar Ray

Institut de Physique Théorique, CEA-Saclay, F-91191 Gif-sur-Yvette Cedex, France

(Received 11 December 2011; published 3 February 2012)

Generating gaugino masses at the leading order has typically been difficult in direct/semidirect gauge mediated supersymmetry breaking models. The Komargodski-Shih (KS) theorem has established that local stability of the supersymmetry breaking vacuum implies a vanishing leading-order gaugino mass in generic renormalizable O'Raifeartaigh models. We relax the condition of renormalizability and investigate the possibility to evade the KS no-go theorem using higher dimensional operators in the Kähler potential and the superpotential. We demonstrate that higher dimensional terms which are polynomial in superfields are not adequate to evade the KS theorem. We narrow down on the possible class of nonpolynomial corrections that can induce unsuppressed gaugino mass in a global supersymmetry breaking vacuum. We find that these models are tantalizingly close to the theories obtained from strongly coupled supersymmetry breaking schemes.

DOI: 10.1103/PhysRevD.85.035003 PACS numbers: 12.60.Jv

I. INTRODUCTION

The realization that generalized O'Raifeartaigh (O'R) models of direct gauge mediated supersymmetry (SUSY) breaking [1] are low energy description of dynamical supersymmetry breaking scenarios from a strongly coupled sector, has been known for some time now. Better understanding of this phenomenon was achieved in [2], which has kindled renewed interest in these models. A typically stubborn problem of these scenarios is the generation of gaugino masses at the leading order, even with explicit tree level R-symmetry breaking, see [3] for a recent review of direct and semidirect gauge mediation models. First pointed out in [4], explicit calculations with all known renormalizable models of direct gauge mediation have shown that cancellations lead to zero gaugino masses at the leading order whereas scalar masses are generally generated at two loop level. Further, the phenomenon of gaugino mass screening [5] prevents gaugino masses being generated at the next order in the messenger loop.[1] This further complicates the possibility to generate sizable gaugino masses in direct gauge mediation models. It was finally realized in [7] that the condition for local stability of the supersymmetry breaking pseudomoduli direction would prevent gaugino masses from being generated at the leading order for general renormalizable models of direct gauge mediation. It was demonstrated that for stable supersymmetry breaking pseudomoduli direction, the determinant of the fermionic mass matrix for the messengers is independent of the pseudomoduli field dependence. This leads to a vanishing gaugino mass in the leading order which is proportional to, $M_g^a \propto \partial \,\mathrm{logdet}(M_f)/\partial X$ where X is the pseudomoduli field and M_f is the fermionic mass matrix for the messenger fields. The vanishing gaugino masses in direct gauge mediation models is now understood in terms of this Komargodski-Shih (KS) no-go theorem.

With the early data from the LHC [8] constraining the SUSY spectra in general and the gluino, in particular, to be relatively heavy, it has become evermore important to investigate avenues to generate unsuppressed gaugino masses in direct gauge mediation models of supersymmetry breaking. Recently, ways to ameliorate this problem have been suggested in the literature [9,10]. In [9] the discussion is based on the fact that the form of the fermionic mass matrix for the messenger fields is not constrained by the KS theorem for models with tachyonic directions in the scalar potential. One would expect leading-order gaugino masses to be generated is these models. Noncanonical Kähler corrections can be used in these models to lift the tachyonic directions. It has been argued that with the noncanonical Kähler corrections, the effective scalar potential of these models will not have any tachyonic direction but leading-order gaugino masses will be generated. In the present paper we make a complementary investigation. We study the possibility to evade the KS theorem by introducing nonrenormalizable terms to models with stable supersymmetry breaking vacuum. We consider the possibility that these contributions introduce a holomorphic pseudomoduli dependence in the determinant of the fermionic mass matrix for the messengers generating leading-order gaugino masses, without disturbing the vacuum configuration.

We investigate the possibility to generate leading-order gaugino masses by introducing nonrenormalizable operators in both the superpotential and the Kähler potential. The most general form of the noncanonical Kähler terms that can contribute to the reduced fermionic mass matrix of the messenger fields are identified. We note that all possible nonrenormalizable superpotential terms can be considered to be a subset of the noncanonical Kähler terms

[1]However see [6] for ways to address this problem by using chiral messengers.

as far as their contribution to the messenger mass matrices in the desired vacuum is considered. We systematically study the viability of generating unsuppressed gaugino masses using higher dimensional terms that are polynomial in the fields. Though we do not specify the UV completion of these models, they can in principle be considered to have originated from some perturbative dynamics at higher energy. However, we find that this class of models are unable to generate unconstrained gaugino masses which are in general suppressed by the high cutoff scale (Λ) of the effective nonrenormalizable theory. The lowest order Kähler term which induces nontrivial corrections to the fermionic mass matrix of the messenger fields has a mass dimension of four. Qualitatively, we observe that beyond this order, gaugino masses are suppressed by $(\langle X\rangle/\Lambda)^{\delta-4}$ where $\langle X\rangle$ is the vev of the pseudomoduli field and δ is the dimension of the operator in the Kähler potential. In general one expects $\langle X\rangle \ll \Lambda$, hence a large suppression.

Next we relax the condition of perturbative UV completion and consider more general functions of the fields motivated by strongly coupled supersymmetry breaking scenarios. We demonstrate that with this generalization the condition for local stability can be explicitly solved in the simplest cases. We obtain surprisingly simple solutions for models of supersymmetry breaking that evade the KS theorem. This class of models break supersymmetry at the global minimum but generates unconstrained gaugino masses. However the condition of local stability of the pseudomoduli direction puts severe constraints on the functional form of the effective Goldstino-messenger terms in the superpotential. The general class of interactions that are allowed are very close to the UV complete theories studied in the literature.

The rest of the paper is organized as follows: In Sec. II, we briefly review the KS theorem within the renormalizable setup and then lay down the framework to generalize to nonrenormalizable scenarios. In Sec. III, we consider the possibility to evade the KS theorem using higher dimensional operators that are polynomial in fields. In Sec. IV, we consider the nonpolynomial generalization. Finally in Sec. V, we conclude with some general observations.

II. GENERALIZATION OF THE KS THEOREM

A. A review of the KS theorem in the renormalizable scenario

Consider a general O'R theory with canonical Kähler potential and a renormalizable superpotential. Let the gauge singlet X and $\{\phi_a\}$ be a set of chiral superfields which constitutes the sector that will break SUSY spontaneously. In order that the $\{\phi_a\}$ should also act as messengers, they should be charged under the SM gauge group. The superfield X is an SM singlet and can get a vev in the vacuum configuration to break SUSY spontaneously. Typically X represents a flat direction in the scalar potential. With this field content, the most general renormalizable superpotential can be written as

$$W = fX + \tfrac{1}{2}(\lambda_{ab}X + m_{ab})\phi_a\phi_b + \tfrac{1}{6}g_{abc}\phi_a\phi_b\phi_c. \quad (1)$$

Here the fermionic mass matrix for the messenger fields $\{\phi_a\}$ is $\mathcal{M}_{\mathcal{F}} = W_{ab} = \lambda_{ab}X + m_{ab}$, where $W_a \equiv \partial W/\partial\phi_a$. In general the determinant of this matrix may be written as

$$\det(\mathcal{M}_{\mathcal{F}}) = \det(\lambda_{ab}X + m_{ab}) = \Sigma C_n(\lambda, m)X^n. \quad (2)$$

Let the roots of the polynomial on the right-hand side of the above equation be defined by $\Sigma C_n(\lambda, m)X^n|_{X\to X_0^l} = 0$. At $X = X_0^l$ the determinant of the fermionic mass matrix vanishes and a Goldstino direction (v) is defined for every root of the polynomial as follows:

$$(\lambda_{ab}X_0^l + m_{ab})v = 0. \quad (3)$$

The bosonic mass matrix for the messenger fields is given by

$$\mathcal{M}_B^2 = \begin{pmatrix} \mathcal{M}_{\mathcal{F}}^*\mathcal{M}_F & \mathcal{F}^* \\ \mathcal{F} & \mathcal{M}_{\mathcal{F}}\mathcal{M}_{\mathcal{F}}^* \end{pmatrix}, \quad (4)$$

where $\mathcal{F}_{ab} = W_c^* W_{abc}$. If the pseudomoduli direction is locally stable everywhere then the scalar mass matrix has to be positive semidefinite. However note that if v is the Goldstino direction defined by $\mathcal{M}_{\mathcal{F}}v = 0$ then it is easy to show that

$$\begin{pmatrix} v \\ v^* \end{pmatrix}^\dagger \begin{pmatrix} M_{\mathcal{F}}^*\mathcal{M}_{\mathcal{F}} & \mathcal{F}^* \\ \mathcal{F} & M_{\mathcal{F}}\mathcal{M}_{\mathcal{F}}^* \end{pmatrix} \begin{pmatrix} v \\ v^* \end{pmatrix} = v^T\mathcal{F}v + cc. \quad (5)$$

The right-hand side of this equation must vanish identically if the bosonic mass matrix is required to be positive semidefinite, otherwise one can make the expression negative by rotating the complex phase of v. We conclude that the condition of local stability of the desired vacuum implies that for a massless Goldstino (v) in the fermionic sector there exists a flat direction in the scalar potential given by the vector (vv^*). An important corollary of this is

$$\mathcal{F}_{ab}v = f\lambda_{ab}v = 0. \quad (6)$$

Using Eq. (6) in Eq. (3) we find v has to be a simultaneous null eigenvector of the matrices λ_{ab} and m_{ab}. This implies that v is a null eigenvector of any matrix of the form $\alpha\lambda_{ab} + \beta m_{ab}$. It follows that $\det(\mathcal{M}_{\mathcal{F}}) = 0$, contradicting our original assumption that the determinant is not identically zero. Thus we find that the assumption taken in Eq. (2) is inconsistent and we conclude that

$$\det(\mathcal{M}_{\mathcal{F}}) = \det(\lambda_{ab}X + m_{ab}) = \text{Const.} \quad (7)$$

It follows that the leading-order gaugino masses given by

$$M_g^a \sim \frac{\alpha^a}{4\pi}\bar{W}_{\bar{X}}\frac{\partial}{\partial X}\,\text{logdet}(M_f), \quad (8)$$

vanish. In conclusion the KS theorem demonstrates that in renormalizable models of direct gauge mediation with a locally stable pseudomoduli direction, gaugino masses are not generated at the leading order.

B. Nonrenormalizable generalization

To study this scenario in the nonrenormalizable setup we first define the desired vacuum configuration of a theory with the field content of Sec. II A. In order to preserve the SM gauge group we should have $\langle\phi_a\rangle = 0 \forall a$. The only field that can take a vev to spontaneously break SUSY is X. Hence we are looking at a vacuum of the form

$$\langle X\rangle \rightarrow \text{undetermined}, \qquad \{\langle\phi_a\rangle = 0\} \quad \forall a. \tag{9}$$

We start with the general renormalizable superpotential given in Eq. (1). The superpotential is linear in X representing a flat pseudomoduli direction in the scalar potential. We find that the two equations $W_X = W_{\phi_i} = 0$ cannot be simultaneously satisfied. At the desired vacuum we have $\langle W_{\phi_i}\rangle = 0, \langle W_X\rangle = f$ and SUSY is broken spontaneously. Considering that the flat direction is locally stable everywhere the determinant of the reduced fermionic mass matrix for the messenger fields remains independent of the pseudomoduli field by the KS theorem implying a zero gaugino mass at the leading order. Our objective is to introduce an X dependence into the determinant of reduced fermionic mass matrix for the messenger fields by adding nonrenormalizable terms to a theory like this without disturbing the local stability of the SUSY breaking vacuum. We will consider nonrenormalizable terms both in the superpotential and in the Kähler potential that can generate such corrections to the mass matrices at the vacuum configuration.

We first consider noncanonical Kähler terms. Following the notations of [11], the messenger mass matrices for the generic noncanonical Kähler potential can be written as

$$\mathcal{M}_{\mathcal{F}}{}^{\mathrm{NC}} = \mathcal{M}_{\mathcal{F}}{}^{C} - \Gamma^d_{ab} W_d, \tag{10}$$

where $\Gamma^d_{ab} W_d = (K^{d\bar{e}}\partial_a K_{b\bar{e}})W_d$. The bosonic mass matrix also receives further corrections due to the noncanonical Kähler terms and can be given as

$$(\mathcal{M}_B^{\mathrm{NC}})^2 = \begin{pmatrix} \mathcal{M}_{\mathcal{F}}{}^{\mathrm{NC}}\mathcal{M}^{*}_{\mathcal{F}}{}^{\mathrm{NC}} - \bar{W}_{\bar{a}}(R_{\bar{b}b})^{a\bar{a}}W_a & \mathcal{F}^{*\mathrm{NC}} \\ \mathcal{F}^{\mathrm{NC}} & \mathcal{M}_{\mathcal{F}}{}^{*\mathrm{NC}}\mathcal{M}_{\mathcal{F}}{}^{\mathrm{NC}} - \bar{W}_{\bar{a}}(R_{\bar{b}b})^{a\bar{a}}W_a \end{pmatrix}, \tag{11}$$

where $\bar{W}_{\bar{a}}(R_{\bar{b}b})^{a\bar{a}}W_a = \bar{W}_{\bar{a}}(K^{\bar{a}c}\partial_{\bar{b}}\Gamma^a_{bc})W_a$ and $\mathcal{F}^{\mathrm{NC}} = \partial_{bc}(W_a K^{\bar{a}a})\bar{W}_{\bar{a}}$. Considering that in the vacuum we can only have $W_X = \bar{W}_{\bar{X}} \neq 0$, the nonzero components are given by, $(R_{\bar{b}b})^{X\bar{X}} \sim K^{\bar{X}c}\partial_{\bar{b}}\Gamma^X_{bc}$ and $\mathcal{F}^{\mathrm{NC}} \sim \partial_{bc}(W_a K^{\bar{X}a})\bar{W}_{\bar{X}}$.

By inspecting Eq. (10), one can see that the new terms need to be bilinear and holomorphic in the messenger fields in order to contribute to the fermionic messenger mass matrices. Thus the most general structure of the noncanonical part of the Kähler potential that contributes to the fermionic mass matrices of the messenger fields may be symbolically represented as

$$K \supset C_{ab}\phi_a\phi_b f\left(\frac{X}{\Lambda}, \frac{\bar{X}}{\Lambda}\right) + cc, \tag{12}$$

where $C_{ab} \neq 0 \Leftrightarrow Q(\phi_a\phi_b) = 0$ and all other terms are zero. $Q(\hat{O})$ represents all the charges of the operator Ô under the SM gauge groups.

With this form of the Kähler terms the curvature tensor $\bar{W}_{\bar{a}}(R_{\bar{b}b})^{a\bar{a}}W_a = 0$. We note that the presence of a nonzero curvature tensor in the Kähler metric results in new contribution to the gaugino masses. With these new contributions it is impossible to recast the scalar and fermionic messenger mass matrices in the form

$$W^{\mathrm{mess}}_{\mathrm{eff}} = M_{ab}\phi_a\tilde{\phi}_b + \theta^2 F_{ab}\phi_a\tilde{\phi}_b,$$
$$\mathcal{L}^{\mathrm{mess}}_{\mathrm{eff}} = -(M_{ab}\psi_a\bar{\psi}_b + \mathrm{H.c.}) - (\varphi_a\tilde{\varphi}^*_a) \times \begin{pmatrix} MM^\dagger & F^* \\ F & M^\dagger M \end{pmatrix}\begin{pmatrix} \varphi^*_b \\ \tilde{\varphi}_b \end{pmatrix}, \tag{13}$$

where ψ and φ are the fermionic and scalar component, respectively, of the chiral messenger superfield ϕ. This would potentially cause the generated gaugino masses to deviate from the expression given in Eq. (8). This in itself is an interesting avenue to generate leading-order gaugino masses in direct gauge mediation models and needs to be explored further. However the arguments of the KS theorem crucially depend on the expression for the gaugino masses as given by Eq. (8) and are not well understood in scenarios where this is no longer true. In this paper we will be confined to models where the curvature tensor identically vanishes. With this choice the only new contributions to the mass matrices are given by

$$\mathcal{M}_{\mathcal{F}}{}^{\mathrm{NC}} = \mathcal{M}_{\mathcal{F}}{}^{C} - C_{ab}\langle W_X\rangle f_X\left(\frac{X}{\Lambda}, \frac{\bar{X}}{\Lambda}\right), \tag{14}$$

$$\mathcal{F}^{\mathrm{NC}} = \mathcal{F}^{C} - C_{ab}|W_X|^2 f_{X\bar{X}}\left(\frac{X}{\Lambda}, \frac{\bar{X}}{\Lambda}\right), \tag{15}$$

where $f_x \equiv \partial f/\partial x$.

At this stage we note that the arguments for the KS theorem used in the canonical case are no longer applicable. We find that if v is now a simultaneous eigenvector of both $\mathcal{F}^{\rm NC}$ and ${\mathcal{M}_{\mathcal{F}}}^{\rm NC}$ one cannot argue that the determinant of ${\mathcal{M}_{\mathcal{F}}}^{\rm NC}$ has to be identically zero everywhere. This is because the matrix form of $\mathcal{F}^{\rm NC}$ is in general different from ${\mathcal{M}_{\mathcal{F}}}^{\rm NC}$. They also have different dependences on X and/or $\bar{X}$. Some generic observations are now in order:

(i) The KS argument is valid only in case of a locally stable pseudomoduli directions i.e., for scenarios where the reduced scalar messenger mass matrix is positive semidefinite. The assertion that new contributions from the noncanonical Kahler potentials can evade this argument and generate leading-order gaugino masses should be supplemented by an example by example demonstration that these additional terms should not destabilize the scalar mass matrix.

(ii) Corrections to the Kähler terms can potentially lead to wrong sign kinetic terms in certain region of the field space. And this consideration puts stringent constraints on the possible form of higher dimensional corrections that are allowed in the Kähler potential. However one can assume that high energy dynamics near the cutoff scale can fix this malady. We will ignore this consideration with the understanding that cutoff scale is much larger than the scale of SUSY breaking.

We now turn our attention to possible nonrenormalizable superpotential terms. The most general superpotential term that contributes to the fermionic mass matrix for the messenger fields, in the vacuum configuration defined in Eq. (9) is given by

$$\Delta W_{\rm NR} = m\phi_a\phi_b g\left(\frac{X}{\Lambda}\right). \tag{16}$$

The contribution of this term to the mass matrices in the desired vacuum configuration is identical to the Kähler potential given in Eq. (12) with the following identifications,

$$f\left(\frac{X}{\Lambda},\frac{\bar{X}}{\Lambda}\right) = \frac{\bar{X}}{\Lambda}g\left(\frac{X}{\Lambda}\right) \quad \text{and} \quad m = \frac{C_{ab}\langle W_X\rangle}{\Lambda}. \tag{17}$$

Thus we note that the most general nonrenormalizable terms that can be added to the superpotential and can contribute to the mass matrices are a specific subset of the most general noncanonical Kähler terms as far as their contribution in the vacuum configuration is considered. It follows that a study of the effect of nonrenormalizable terms in direct gauge mediation models can be effectively carried out by considering the noncanonical terms in the Kähler potential alone.

Having made this observation it should be noted that there are definite differences between a higher dimensional superpotential term and a noncanonical Kähler term. These differences show up in the global structure of the scalar potential specifically in the field space regions away from the SUSY breaking vacuum.

III. THEORIES WITH POLYNOMIAL CORRECTIONS

If we consider a perturbative UV completion of the theories, we can expect these effective terms to be generated by integrating out heavy states operative at high scale. This consideration constraints the functional form of f defined in Eqs. (12) and (16) to be a polynomial of the fields. In this section we will discuss the possibility of evading the KS theorem to generate unconstrained gaugino masses using such polynomial correction to the Kähler potential and the superpotential.

A. Noncanonical Kähler potentials:

Let us consider that the function f in Eq. (12) is a polynomial in both X and $\bar{X}$. Thus generically we may write,

$$f\left(\frac{X}{\Lambda},\frac{\bar{X}}{\Lambda}\right) = \sum_{n\bar{n}} C^{n\bar{n}}\frac{X^n\bar{X}^{\bar{n}}}{\Lambda^{n+\bar{n}}}. \tag{18}$$

In this case the contributions to the matrices are of the following form:

$${\mathcal{M}_{\mathcal{F}}}^{\rm NC} = {\mathcal{M}_{\mathcal{F}}}^{C} - \sum_{n\bar{n}} C^{n\bar{n}}_{ab}\langle W_X\rangle \bar{n}\frac{X^n\bar{X}^{\bar{n}-1}}{\Lambda^{n+\bar{n}}}, \tag{19}$$

$$\mathcal{F}^{\rm NC} = \mathcal{F}^{C} - \sum_{n\bar{n}} C^{n\bar{n}}_{ab}|W_X|^2\bar{n}n\frac{X^{n-1}\bar{X}^{\bar{n}-1}}{\Lambda^{n+\bar{n}}}. \tag{20}$$

It is clear from Eqs. (19) and (20) that for the new nonrenormalizable terms to contribute we should ensure $\bar{n}\neq 0$. We will now summarize how the individual terms contribute to the gaugino mass and the stability condition for various choices of $n, \bar{n}$.

(i) The lowest order contribution comes from the term $\bar{n}=1, n=0$ In this case we find that the new contribution is just a redefinition of the matrix $m_{ab}\to m_{ab} - C_{ab}\langle W_X\rangle/\Lambda$. We can now trace the arguments given in Sec. II A identically. This will naturally lead to the conclusion that if the vacuum is locally stable, leading-order gaugino masses will vanish.

(ii) The next order contribution comes when $n=1, \bar{n}=1$. In this case we find that the contribution simply results in a redefinition of the matrix $\lambda_{ab}\to \lambda_{ab} - C_{ab}\langle W_X\rangle/\Lambda^2$. This again leads to the same conclusion as in the previous case.

(iii) At this same order we have a nontrivial contribution given by $\bar{n}=2, n=0$. This contributes to the fermionic mass matrix but does not contribute to $\mathcal{F}$. This cannot be modeled by redefinition of parameters. However we make the observation that this term cannot directly introduce a holomorphic dependence on X, into the fermionic mass matrix. With the observation that $\det(\lambda_{ab}X + m_{ab}) =$ Const, we expect the $\det(\lambda_{ab}X + m_{ab} - C_{ab}\langle W_X\rangle\bar{X}/\Lambda^2)\sim \bar{X}X/\Lambda^2$. This will lead to gaugino mass terms that are suppressed by the factor $\langle \bar{X} \rangle / \Lambda$. In general it is well known that in O'R models the one loop correction fixes the X vev near zero [12]. This will certainly be modified due to the presence of the noncanonical Kähler terms. It is still expected that the vev will be generally at a scale where $\langle X \rangle \ll \Lambda$ and thus lead to a suppression of the generated gaugino masses.

(iv) All higher order noncanonical Kähler terms with $\bar{n} + n > 2$ will in general lead to further suppression in the gaugino mass terms of the order $(\frac{\langle X \rangle}{\Lambda})^{n-1} (\frac{\langle \bar{X} \rangle}{\Lambda})^{\bar{n}-1}$.

In conclusion we observe the generic noncanonical Kähler terms of perturbative origin when added to O'R models with global SUSY breaking can only lead to leading-order gaugino masses which are suppressed by the cutoff scale. This general observation is made without any reference to the stability condition of the vacuum. Note that in this class of models the determinant of the fermionic mass matrix will be a polynomial in X and therefore will have roots in the finite complex plane. The pseudomoduli direction will in general have an instability at the point where the determinant vanishes.

B. Nonrenormalizable superpotential terms:

In continuation of the discussion in the previous section we point out that the most general nonrenormalizable terms in the superpotential which are polynomial in the superfields are a subset of the Kähler potential defined in Eq. (18). In the phenomenologically acceptable vacuum, the contribution to the messenger mass matrices from these noncanonical Kähler terms with $\bar{n} = 1$ corresponds to the contribution from the most general nonrenormalizable superpotential term given by

$$\Delta W = \sum_n m_{ab}^{(n)} \phi_a \phi_b \left(\frac{X}{\Lambda} \right)^n, \tag{21}$$

where $m^{(n)}$ can be read off from Eq. (17). The limitations of such terms for $n = 0, 1, >1$ are similar to the ones discussed earlier.

We make the general observation that starting with a direct gauge mediation theory where SUSY is broken globally and the leading-order gaugino masses disappear due to the KS theorem, it is impossible to generate them by adding nonrenormalizable terms that are polynomial in the fields, either to the superpotential or the Kähler potential.

IV. THEORIES WITH NONPOLYNOMIAL CORRECTION

With the conclusion of the previous section we abandon the possibility of circumventing the KS theorem using higher dimensional terms that are polynomial in the superfields, possibly arising from perturbative dynamics at high energy scales. Instead we turn our attention to terms arising from theories with nonperturbative UV completion. Effective low energy description of nonperturbative theories of SUSY breaking can give rise to terms that are nonpolynomial in the superfields. The theories of dynamical SUSY breaking [13,14], commonly incorporate terms that are exponential of the superfields. In theories where gaugino condensates are utilized to break SUSY, the exponential of the dilaton fields commonly appears [15]. In retrofitted O'R models [16] where the vev of the pseudomoduli is dynamically generated, we find the effective superpotential at energies below the dynamical scale contains terms where the pseudomoduli superfields appear in the exponential. Nonpolynomial terms arise in the effective superpotential of SUSY theories with Intriligator, Seiberg and Shih type supersymmetry breaking. This is essentially generated from the dual of nonperturbative strongly coupled supersymmetric quantum chromodynamics like theories [17,18]. In this class of theories the pseudomoduli field commonly appears with negative powers in the superpotential and the Kähler potential. In the present paper our paradigm is to take a bottom up approach to the problem of generating leading-order gaugino masses in the O'R models, thus evading the KS theorem. We will neither endeavor to construct a UV complete theory of the hidden sector nor try to demonstrate the ability to evade the KS theorem with nonpolynomial terms in complete generality. Rather our approach will be to investigate this as a possibility using examples.

To keep matters simple we will look at the possibility of adding a nonrenormalizable superpotential term to theories that break supersymmetry globally. We will consider the simplest supersymmetry breaking scenario. Let X be the Standard Model gauge singlet chiral superfield. And $(\phi \tilde{\phi})$ is a vectorlike[2] pair of messenger fields charged under the standard model gauge group. The simplest SUSY breaking sector that can be constructed with this field content is given by the following superpotential:

$$W = -\mu^2 X + f(X) \phi \tilde{\phi}. \tag{22}$$

We will assume the that the Kähler potential is canonical. The condition that the theory generates nonzero gaugino mass at leading order means that $f(X)$ has to be a nonconstant function of X. If we further demand that the theory breaks supersymmetry globally, one needs to impose the condition $f(X) \neq 0$ everywhere in the finite complex plane. Note that this condition is far stronger than the requirement of local stability which is enough to discuss the KS theorem.

If we insist that the superpotential is holomorphic in the entire complex plane then $f(X)$ should also be an analytic

[2]These charged messenger superfields can be considered to fill a complete representation of a grand unified theory gauge group like the SU(5) required to preserve gauge coupling unification.

function of X. This implies that $f(X)$ is an entire function and subject to the constraints of the little Picard theorem. The examples of entire functions that do not take the value of zero in the entire finite complex plane are limited. From a phenomenological perspective a well-motivated choice would be to take $f(X) = me^{-(X/\Lambda)}$ in Eq. (22). This is the simplest entire function that is nonzero everywhere in the finite complex plane. Thus we expect SUSY to be broken globally in this model. In the desired vacuum the mass matrices for the messenger fields now take the following form:

$$m_f = me^{-(X/\Lambda)} \quad \text{and}$$
$$m_B^2 = \begin{pmatrix} m^2 e^{-(X+X^*/\Lambda)} & \frac{m\mu^2}{\Lambda} e^{-(X^*/\Lambda)} \\ \frac{m\mu^2}{\Lambda} e^{-(X/\Lambda)} & m^2 e^{-(X+X^*/\Lambda)} \end{pmatrix}. \tag{23}$$

The condition for local stability of the pseudomoduli direction now reduces to

$$|m^2 e^{-(X+X^*/\Lambda)}| < |\frac{m\mu^2}{\Lambda} e^{-(X/\Lambda)}|. \tag{24}$$

As is evident, this condition is easily violated at finite values of X, rendering the vacuum unstable at that point. Typically, these instabilities leads to a vacuum with anomalous breaking of the Standard Model gauge group. It should be noted that this conclusion is not an artifact of the simple form of the superpotential considered and it cannot be resolved by a simple enlargement of the messenger sector.

A. Generic solution to the local stability condition

Finally, we abandon the constraint that $f(X)$ is analytic everywhere. Rather we directly try to solve for condition of local stability. Using Eq. (22), the scalar mass matrix for the messenger fields is given by

$$m_B^2 = \begin{pmatrix} |f(X)|^2 & -\left(\mu^2 \frac{\partial f(X)}{\partial X}\right)^* \\ -\mu^2 \frac{\partial f(X)}{\partial X} & |f(X)|^2 \end{pmatrix}. \tag{25}$$

To establish that a 2×2 matrix is positive definite it is enough to show that the trace and the determinant are positive. The condition on the trace is trivially satisfied by the above matrix. We turn our attention to the determinant. The condition that the determinant has to be positive implies

$$|f(X)|^4 \geq \left| \mu^2 \frac{\partial f(X)}{\partial X} \right|^2. \tag{26}$$

We consider the scenario that saturates this bound. To solve the resulting equation we separate the real and the complex parts, giving the relation

$$\frac{f(X)^2}{\mu^2 \partial f(X)/\partial X} = \left(\frac{f(X)^2}{\mu^2 \partial f(X)/\partial X}\right)^* = e^{i\theta}. \tag{27}$$

This simplifies to the following differential equation:

$$f(X)^2 = e^{i\theta} \mu^2 \frac{\partial f(X)}{\partial X}. \tag{28}$$

The functional form of $f(X)$ can be easily obtained by solving the differential equation which gives us

$$f(X) = \frac{\mu^2 e^{i\theta}}{X + b}. \tag{29}$$

Note that this solution saturates the bound given in Eq. (26). Without any loss of generality we can choose the function to be $f(X) = m^2/X$, where m is a real constant. We observe $f(X)$ though not defined at $X = 0$, is analytic everywhere else. As long as $\langle X \rangle \neq 0$, the theory defined by the superpotential given in Eq. (22) is well behaved. To demonstrate the local stability of this theory we consider the scalar mass matrix which now takes the following form:

$$m_B^2 = \begin{pmatrix} \frac{m^4}{|X|^2} & \frac{m^2\mu^2}{(X^*)^2} \\ \frac{m^2\mu^2}{X^2} & \frac{m^4}{|X|^2} \end{pmatrix}. \tag{30}$$

We note that the eigenvalues of this matrix are given by $(m^2 - \mu^2)m^2/|X|^2$ and $(m^2 + \mu^2)m^2/|X|^2$. Thus, for $m^2 > \mu^2$, the eigenvalues are positive for any value of $\langle X \rangle$ and matrix is positive definite. Therefore with this constraint on the parameters the pseudomoduli direction is locally stable everywhere. Importantly, we also note that $f(X)$ does not take the value zero in the finite complex plane. This means that not only the pseudomoduli direction is locally stable everywhere, supersymmetry is also broken globally. It naturally satisfies all the conditions we laid down on $f(X)$ at the beginning of this section. Let us now investigate the global structure of the scalar potential. The potential $V = \sum_a W_a$ where,

$$W_X = -\mu^2 - m^2\phi\tilde{\phi}/X^2, \qquad W_\phi = m^2\tilde{\phi}/X,$$
$$W_{\tilde{\phi}} = m^2\phi/X. \tag{31}$$

Clearly these three equations cannot be simultaneously put to zero and supersymmetry is broken globally. Curiously the condition $m^2 > \mu^2$ implies that there is only one global minimum[3] of the potential given by $\langle X \rangle \rightarrow$ undetermined and $\langle \phi \rangle = \langle \tilde{\phi} \rangle = 0$ and $V = \mu^4$. With the single constraint on the superpotential parameters, we not only ensure that the desired vacuum is locally stable but also enforce it to be the global minimum of the scalar potential.

[3]A lower lying minimum only appears when $m^2 < \mu^2$, in this case the minimum is at $V = (\mu^2 - m^2)m^2$.

The fermionic mass matrix for the messenger is simply given by

$$\det(m_f) = m^2/X. \tag{32}$$

Gaugino masses are generated at the leading order. Using Eq. (8) and Eq. (32) we find that

$$M_a \sim \frac{\alpha_a}{4\pi}\mu^2\frac{1}{\langle X\rangle}, \tag{33}$$

which is unsuppressed by any high scale. And unlike the minimal gauge mediation models, within this framework the messenger masses may be in the TeV scale and observable at the present collider experiments. This will potentially lead to interesting phenomenological scenarios at collider experiments.

In conclusion we note that the possibility to generate gaugino masses at leading order through direct gauge mediation with locally stable SUSY breaking vacuum is restricted to very specific class of models even in its nonperturbative generalization. Crucially the interactions of the pseudomoduli field with the messengers are restricted to have very specific functional forms. This brings us to the possible origin of this class of superpotentials. It is well known that models of supersymmetry breaking with an supersymmetric quantum chromodynamics sector generate effective superpotentials at low energies which have the pseudomoduli fields appearing in the denominator [18]. However, we could not find an instance in the literature where the effective term discussed here appears in its exact form. To the best of our knowledge, such terms can not be generated within the framework of the simplest nonperturbative scenarios like the Intriligator, Seiberg and Shih .

V. CONCLUSION

In this paper we have studied the possibility of adding simple nonrenormalizable terms to globally stable SUSY breaking O'R models to evade the KS no-go theorem. This is complementary to the study carried out in [9] where unstable renormalizable theories were considered and noncanonical Kähler terms were used to lift these instabilities.

Within this framework we have demonstrated that the simple higher dimensional terms which are polynomial in the fields, and thus can potentially be generated through perturbative dynamics at higher scales, are not adequate to alleviate the problem of generating large unconstrained gaugino masses. Typically we find in these models the gaugino masses are suppressed by the high cutoff scale of the effective theory. Further they exhibit tachyonic directions along the pseudomoduli direction at points where the determinant of the fermionic mass matrix vanishes.

Next we have considered nonpolynomial terms that can generate unconstrained gaugino masses without disturbing the stability of the vacuum. In this context we have imposed a stronger constraint on the theory, demanding that the desired SUSY breaking vacuum is the global minimum of the scalar potential. With these restrictive constraints we solved for the condition of local stability of the potential. We obtain a surprisingly simple solution that satisfies all the conditions of local and global stability and generates unsuppressed gaugino masses at the leading order. We observe that supersymmetry breaking models having these virtues will have a very specific form of superpotential where the pseudomoduli field couples to messenger field with inverse one power. This might have consequences for Goldstino couplings and can have major cosmological impact. A systematic discussion of these issues is beyond the mandate of this paper. The form of the nonpolynomial terms required for this is also tantalizingly close to the ones that originate from generic nonperturbative schemes of SUSY breaking discussed in the literature.

A more thorough study of possible nonpolynomial terms described in the literature should be carried out in the context of direct gauge mediation models. The possibility of using them to evade the KS theorem and generate phenomenologically viable soft SUSY breaking spectrum needs to be carried out. In this context we also note that the entire discussion in this paper is carried out within a framework where the Kähler metric is flat i.e., the curvature tensor is considered to be zero everywhere. Relaxation of this constraint may lead to more phenomenologically acceptable avenues to evade the KS theorem.

ACKNOWLEDGMENTS

This work would not have been possible without the support and guidance of S. Lavignac. The work of T. S. R. is supported by the EU ITN under Contract "UNILHC" No. PITN-GA-2009-237920, the CEA-Eurotalent program and the Agence Nationale de la Recherche under Contract No. ANR 2010 BLANC 0413 01.

[1] G.F. Giudice and R. Rattazzi, Phys. Rep. 322, 419 (1999).
[2] K. A. Intriligator, N. Seiberg, and D. Shih, J. High Energy Phys. 04 (2006) 021.K. I. Izawa and T. Yanagida, Prog. Theor. Phys. 95, 829 (1996).K. A. Intriligator and S. D. Thomas, Nucl. Phys. B473, 121 (1996).
[3] R. Kitano, H. Ooguri, and Y. Ookouchi, arXiv:1001.4535.
[4] J. Polchinski and L. Susskind, Phys. Rev. D 26, 3661 (1982).

[5] N. Arkani-Hamed, G.F. Giudice, M.A. Luty, and R. Rattazzi, Phys. Rev. D **58**, 115005 (1998).
[6] R. Argurio, M. Bertolini, G. Ferretti, and A. Mariotti, J. High Energy Phys. 12 (2010) 064.
[7] Z. Komargodski and D. Shih, J. High Energy Phys. 04 (2009) 093.
[8] S. Chatrchyan *et al.* (CMS Collaboration), Phys. Rev. Lett. **107**, 221804 (2011); S. Chatrchyan *et al.* (CMS Collaboration), J. High Energy Phys. 08 (2011) 155.
[9] Y. Nakai and Y. Ookouchi, J. High Energy Phys. 01 (2011) 093.
[10] N. Maru, Phys. Rev. D **82**, 075015 (2010).
[11] L. G. Aldrovandi and D. Marques, J. High Energy Phys. 05 (2008) 022.
[12] E. Dudas, S. Lavignac, and J. Parmentier, Phys. Lett. B **698**, 162 (2011).
[13] E. Witten, Nucl. Phys. **B188**, 513 (1981).
[14] M. Dine and W. Fischler, Phys. Lett. **110B**, 227 (1982).
[15] M. Dine and J.D. Mason, Rep. Prog. Phys. **74**, 056201 (2011).
[16] M. Dine, J.L. Feng, and E. Silverstein, Phys. Rev. D **74**, 095012 (2006).
[17] N. Seiberg, T. Volansky, and B. Wecht, J. High Energy Phys. 11 (2008) 004.
[18] K.A. Intriligator and N. Seiberg, Classical Quantum Gravity **24**, S741 (2007).Y. Shadmi, arXiv:hep-th/0601076.