Sizable electron/neutron electric dipole moment in $D3/D7$ μ-split supersymmetry

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Within the framework of $\mathcal{N}=1$ gauged supergravity, using a phenomenological model that can be obtained locally as a Swiss-cheese Calabi-Yau string-theoretic compactification with a mobile $D3$-brane localized on a nearly special Lagrangian three cycle in the Calabi-Yau and fluxed stacks of wrapped $D7$-branes, and which provides a natural realization of μ-split supersymmetry (SUSY), we show that in addition to getting a significant value of an [electron/neutron (e/n)] electron dipole moment (EDM) at two-loop level, one can obtain a sizable contribution of (e/n) EDM even at one-loop level due to the presence of heavy supersymmetric fermions nearly isospectral with heavy sfermions. Unlike traditional split SUSY models in which the one-loop diagrams do not give significant contribution to the EDM of the electron/neutron because of very heavy sfermions existing as propagators in the loop, we show that one obtains a "healthy" value of the EDM in our model because of the presence of a heavy Higgsino, neutralino/chargino, and gaugino as fermionic propagators in the loops. The independent CP-violating phases are generated from nontrivial distinct phase factors associated with four Wilson line moduli [identified with first-generation leptons and quarks and their $SU(2)_L$-singlet cousins] as well as the $D3$-brane position moduli (identified with two Higgses), and the same are sufficient to produce overall distinct phase factors corresponding to all possible effective Yukawas as well as effective gauge couplings that we discuss in the context of $\mathcal{N}=1$ gauged supergravity action. However, the complex phases responsible to generate a nonzero EDM at one-loop level mainly appear from an off-diagonal contribution of sfermion as well as Higgs mass matrices at the electroweak scale (EW). In our analysis, we obtain a dominant contribution of the electron/neutron EDM around $d_e/e \equiv \mathcal{O}(10^{-29})$ cm from two-loop diagrams involving heavy sfermions and a light Higgs, and $d_e/e \equiv \mathcal{O}(10^{-32})$ cm from a one-loop diagram involving a heavy chargino and a light Higgs as propagators in the loop. The neutron EDM gets a dominant contribution of the order $d_n/e \equiv \mathcal{O}(10^{-33})$ cm from the one-loop diagram involving SM-like quarks and Higgs. To justify the possibility of obtaining a large EDM value in the case of a Barr-Zee diagram which involves $W^\pm$ and the Higgs (responsible to generate the nontrivial CP-violating phase) in the two-loop diagrams as discussed by Leigh *et al.* [Nucl. Phys. B267, 509 (1986)], we provide an analysis of the same in the context of our $D3/D7$ μ-split SUSY model at the EW scale. By conjecturing that the CP-violating phase can appear from the diagonalization of the Higgs mass matrix obtained in the context of μ-split SUSY, we also get an EDM of the electron/neutron around $\mathcal{O}(10^{-27})$ e cm in the case of the two-loop diagram involving $W^\pm$ bosons.

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I. INTRODUCTION

For the past few decades, string-theoretic models have been considered to provide an excellent framework for possible unification of gravity with all other fundamental forces. To study the phenomenological implications of these models, the same must invoke a particular supersymmetry (SUSY) breaking mechanism (along with the SUSY breaking scale). The phenomenological models mainly rely on the $\mathcal{O}(\mathrm{TeV})$ SUSY breaking scale because this helps solve serious gauge-hierarchy problems, which, in fact, have been considered as a primary motivation to introduce SUSY. However, low-scale SUSY models give rise to many unwanted phenomenological problems, such as flavor-changing neutral currents. Motivated by obtaining an extremely small cosmological constant and the string landscape scenario, an alternative to these assumptions was proposed by Arkani-Hamed and Dimopoulos (dubbed as "split SUSY") in [1] according to which SUSY is broken at an energy scale way beyond the collider search and could be even near the scale of grand unification. The scenario is emerging to be quite interesting from the point of view of phenomenology because of the fact that heavy scalars mostly appearing as virtual particles in most of the particle decay studies help resolve many diverse issues of both particle physics and cosmology. The μ-split SUSY model was proposed in [2] to alleviate the famous μ problem by further splitting the split SUSY by raising the μ paramter to a large value. Though the exact signatures may not be

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foreseeable in the near future via precise measurements to be carried out at the Large Hadron Collider (LHC), indirect methods can be made available to test some of the signatures of this scenario. In this context, the electron dipole moment (EDM) of the electron/neutron serves as another testing ground for the split SUSY scenario. Recently, the ACME Collaboration has reported a new experimental upper limit of $|d_e| < 8.7 \times 10^{-29}\, e\,\mathrm{cm}$ [3], which is an order-of-magnitude improvement in sensitivity as compared to previous limits [4–7]. The current experimental limit on the neutron EDM [8,9] is $|d_n/e| < 0.29 \times 10^{-25}$ cm.

In the Standard Model (SM), the CP-odd phases generated through the Cabibo-Kobayashi-Masakawa (CKM) matrix give a theoretical bound on the EDM which is far below the experimental limits. However, new CP-violating phases can appear in supersymmetric theory models from complex soft SUSY breaking parameters. In addition to this, in string-inspired models, the CP-violating phases are associated with complex Yukawa couplings originating from string compactifications [10–13]. These CP-violating phases associated with complex soft SUSY breaking parameters as well as Yukawa couplings appearing in different supersymmetric models are typically large, i.e., $\mathcal{O}(1)$ and, hence, do not satisfy the current experimental bounds on the electron and on the neutron EDM. One, hence, has to put stringent constraints on the supersymmetric and, in particular, supergravity (SUGRA) models. More specifically, the limits can be satisfied if one considers (i) unnaturally small CP-violating phases of $\mathcal{O}(10^{-2}$–$10^{-3})$, (ii) multi-TeV superpartners in the model, or (iii) internal cancellations between different supersymmetric contributions to the EDM at loop levels. The constraints on the CP-violating phases in the supersymmetric models have been discussed in [14–16], and the systematic analysis of the EDM up to two loops in the context of the minimal supersymmetric Standard Model (MSSM) is provided in [17–20]. In minimal supergravity (mSUGRA) models discussed in the literature [21–26], the EDM's bounds have been reconciled with the experimental limits by showing sufficient cancellations among different supersymmetric contributions without taking into account $\mathcal{O}(>\mathrm{TeV})$ superpartners and any fine-tuning in phase angles. The main difficulty in choosing multi-TeV scalars as an appropriate mechanism to generate the EDM is because the same abandons naturalness and also requires severe fine-tuning while satisfying radiative electroweak (EW) symmetry breaking. However, the nonobservation of sparticles at the LHC may point toward some sort of fine-tuned natural SUSY [27,28] or the high SUSY scale/split SUSY models [29–31]. Therefore, it is interesting to probe high-scale SUSY models, in particular, μ-split SUSY models, to explain the EDM within the reach of experimental limits because the same also helps satisfy the radiative EW symmetry breaking condition by choosing a natural value of μ, hence, alleviating the μ problem. The region of parameter space satisfying the EDM value of the order of experimental limits has been analyzed in [32] in the presence of $\mathcal{O}(\mathrm{TeV})$ superpartners in the mSUGRA model by considering moderate fine-tuning in $\tan\beta$.

Our approach is quite different in that the SUGRA models discussed in the literature, even in the framework of string compactifications, do not rely on the high supersymmetry breaking scale. On the other hand, the typical split supersymmetry models used to study the EDM of the electron/neutron include heavy sfermions but light gauginos and Higgsinos [33,34]. We analyze the EDM of the electron and neutron in the supergravity limit of local large volume $D3/D7$ type IIB compactifications, which provides, to our knowledge, the first realization of the μ-split SUSY scenario (with large gaugino masses). In typical split SUSY models, all possible one-loop contributions to the EDM are highly suppressed by the superheavy scalar masses in the loop, and leading contributions to the EDM start at the two-loop level due to the presence of SM particles and EW charginos and neutralinos in the loops (for the analysis of two-loop Barr-Zee diagrams in different models, see [33–36] and references therein). However, in our model, the gaugino and neutralino/chargino are almost as heavy as neutral scalars except one light Higgs. Based on that, one cannot ignore the contribution of one-loop diagrams because of the partial compensation of the suppression factors appearing from heavy sfermion masses, by heavy fermions' (neutralino, chargino, and gaugino) masses.

Therefore, in this paper, we perform a quantitative analysis of the neutron and electron EDMs for all possible one-loop as well as two-loop diagrams in the context of large volume $D3/D7$ μ-split supersymmetry. The nonzero imaginary phases that appear through mixing between L-hand (left-handed) and R-hand (right-handed) sfermions (sfermions corresponding to the left- and right-handed components of fermions) at the electroweak scale, play an important role. In addition to discussing the one-loop diagrams that exhibit nonzero phases through mixing between sfermions, we also take into account the loop diagrams in which a unique phase appears through mixing between two Higgses at the electroweak scale. In the large volume μ-split SUSY model of [37,38], we have already calculated the eigenvalues of the Higgs mass matrix at the electroweak scale, which, with some fine-tuning, eventually leads to one light Higgs and one heavy Higgs. In this paper, we append the details of the complex phase associated with off-diagonal components of the Higgs mass matrix, too. Because of the presence of a light and a heavy Higgs in our model, one can expect to get a reasonable order of magnitude of the EDM of the electron/neutron from one-loop diagrams involving Higgs and other SM/supersymmetric particles. The complete analysis has been carried out by including other interesting one-loop diagrams which involve sGoldstinos [identified locally with a "big" divisor (bulk) volume modulus in our setup] as scalar particles in the loop. For two-loop diagrams, we mainly focus on the Barr-Zee diagrams which involve a fermion, sfermions, and $W^{\pm}$ as part of an internal loop and are mediated through $h\gamma$ exchange, except one R-parity violating diagram which involves fermions in the internal loop and is mediated through $\nu_L\gamma$ exchange. For the complete analysis, we also calculate the contribution of rainbow-type two-loop diagrams involving R-parity violating as well as R-parity conserving vertices. For all two-loop diagrams discussed in this paper, the complex effective Yukawa couplings [associated with the $e^{\frac{K}{2}}(\mathcal{D}DW)\bar{\chi}\chi$ term in the $\mathcal{N}=1$ gauged supergravity action of [39]] are sufficient to produce nonzero complex phases to generate a nonzero EDM.

The plan for the rest of the paper is as follows. In Sec. II, we elaborate upon our large volume $D3/D7$ model discussed in [38]. We discuss the details of our phenomenological model in Sec. II A and show the same to be realizable locally as the large volume limit of a type IIB Swiss-cheese Calabi-Yau orientifold involving a mobile space-time filling $D3$-brane localized at a nearly special Lagrangian three cycle embedded in the big divisor (hence, the local nature of the model's realization) and multiple fluxed stacks of space-time filling $D7$-branes wrapping the same big divisor in Sec. II B. After providing the geometrical framework of the model in Sec. II B, we briefly mention the phenomenological results that describe the possible identification of Wilson line moduli with first-generation leptons and quarks as well as their $SU(2)_L$-singlet cousins, and $D3$-brane position moduli with two Higgses. Thereafter, we briefly summarize the calculation and results corresponding to the values of soft SUSY breaking parameters as well as the supersymmetric fermionic masses. In Sec. III, we explain the origin of nonzero complex phases obtained in the context of the $\mathcal{N}=1$ gauged supergravity limit of our local $D3/D7$ model. We also argue that phases of effective Yukawa couplings do not change under a renormalization group flow from string scale down to the electroweak scale in our model. In Sec. IV, we turn towards order-of-magnitude estimates of the EDM of the electron/neutron for various possible one-loop diagrams. The effective vertices are calculated by considering the $\mathcal{N}=1$ gauged supergravity action of [39,40]. The complex phases, as already explained, can be made to appear through the complex off-diagonal components of an sfermion/Higgs mass matrix and complex effective Yukawa couplings appearing in all one-loop diagrams. We assume the phases of both off-diagonal components of the scalar mass matrix as well as possible effective Yukawa's to lie in the range $(0,\frac{\pi}{2}]$ in all the calculations. The section has been divided into three subsections. In Sec. IV A, we give a detailed discussion of one-loop diagrams which involve sfermions as scalar propagators and gauginos, neutralinos and SM-like fermions as fermionic propagators, respectively. Higgs doublets as scalar propagators and chargino and SM-like fermions as fermionic propagators, respectively. Here, the nonzero imaginary phases appear through mixing between two Higgses at the electroweak scale in the Higgs mass matrix. In Sec. IV C, we evaluate the contribution of a heavy gravitino and sGoldstino multiplet to the EDM of the electron/neutron. Though the loop diagrams involving the same are divergent, we pick out the finite contributions for the purpose of obtaining an estimate of the EDM of the electron/neutron in the case of a heavy gravitino. In Sec. V, we consider two-loop Barr-Zee diagrams. The section has been divided into three subsections. In Secs. V A and V B, we compute the two-loop diagrams which involve an internal fermion loop and an internal sfermion loop. These diagrams are mediated by γh and $\gamma\nu_L$ exchange. In Sec. V C, we carry out an analysis of two-loop diagrams involving a W-boson loop in our μ-split SUSY model. In Sec. V D, we discuss two-loop rainbow-type diagrams. Section VI has the summary of our results and a discussion. In the Appendix, we evaluate the chargino mass matrix using the $\mathcal{N}=1$ gauged supergravity action in the context of the large volume $D3/D7$ μ-split SUSY setup.

II. THE SETUP

In [38], within the context of type IIB string theory with a space-time filling $D3$-brane and fluxed stacks of $D7$-branes wrapping a divisor along with ED3/ED1-instanton-generated superpotential and world-sheet instanton-corrected Kähler potential, we worked locally close to a nearly special Lagrangian three cycle (6) within a Swiss-cheese-type Calabi-Yau orientifold (various aspects of this setup will be summarized in Sec. II B). But before we do the same, we will first briefly describe in Sec. II A, a model that could be locally realized as a large volume $D3/D7$ Swiss-cheese setup of [38]. In other words, Sec. II A embeds the local model of [38] into a phenomenological model, something which was not done in [38]. In other words, the phonomenological supergravity model discussed in Sec. II A can be locally geometrically engineered via the construct of [38].

A. The model

For an $\mathcal{N}=1$ compactification, we will take the phenomenological Kähler potential of our model to be

$$K_{\rm pheno} = -\ln[-i(\tau-\bar{\tau})] - \ln\left(-i\int_{CY_3}\Omega\wedge\bar{\Omega}\right) - 2\ln\left[a_B(\sigma_B+\bar{\sigma}_B-\gamma K_{\rm geom})^{\frac{3}{2}} - \left(\sum_i a_{S,i}(\sigma_{S,i}+\bar{\sigma}_{S,i}-\gamma K_{\rm geom})\right)^{\frac{3}{2}} + \mathcal{O}(1)\mathcal{V}\right], \quad (1)$$

where the divisor volumes σ_α are expressible in terms of "Kähler" coordinates T_α, $\mathcal{M}_\mathcal{I}$,

$$\sigma_\alpha \sim T_\alpha - [i\mathcal{K}_{\alpha bc}c^b\mathcal{B}^c + iC_\alpha^{\mathcal{M}_\mathcal{I}\bar{\mathcal{M}}_{\bar{\mathcal{J}}}}(\mathcal{V})\mathrm{Tr}(\mathcal{M}_\mathcal{I}\mathcal{M}^\dagger_{\bar{\mathcal{J}}})], \quad (2)$$

$\alpha = (B, \{S, i\})$ and $\mathcal{M}_\mathcal{I}$ being $SU(3_c) \times SU(2)_L$ bifundamental matter field $a_{\mathcal{I}=2}$, $SU(3_c) \times U(1)_R$ bifundamental matter field $a_{\mathcal{I}=4}$, $SU(2)_L \times U(1)_L$ bifundamental matter field $a_{\mathcal{I}=1}$, $U(1)_L \times U(1)_R$ bifundamental matter field $a_{\mathcal{I}=3}$ along with $SU(2)_L \times U(1)_L$ bifundamental $\tilde{z}_{1,2}$ with the intersection matrix $C_\alpha^{a_I\bar{a}_{\bar{J}}} \sim \delta_\alpha^B C_\alpha^{I\bar{J}}$, $C_\alpha^{a_I\bar{\tilde{z}}_{\bar{j}}} = 0$, $\rho_{S,B}$, $\mathcal{G}^a = c^a - \tau b^a$ being complex axionic fields (α, a running over the real dimensionality of a subspace of the internal manifold's cohomology complex), and the phenomenological superpotential is given as under

$$W_{\mathrm{pheno}} \sim (z_1^{18} + z_2^{18})^{n^s} e^{-n^s \mathrm{vol}(\Sigma_S) - (\alpha_S z_1^2 + \beta_S z_2^2 + \gamma_S z_1 z_2)}, \quad (3)$$

where the bifundamental $\tilde{z}_i$ in K will be equivalent to the $z_{1,2} \in \mathbb{C}$ in W. It is expected that $\mathcal{M}_\mathcal{I}$, $T_{S,B}$, $\mathcal{G}^a$ will constitute the $\mathcal{N} = 1$ chiral coordinates. The intersection matrix elements $\kappa_{S/Bab}$ and the volume-dependent $C_\alpha^{\mathcal{M}_\mathcal{I}\bar{\mathcal{M}}_{\bar{\mathcal{J}}}}(\mathcal{V})$ are chosen in such a way that at a local (metastable) minimum,

$$\begin{aligned}\langle\sigma_S\rangle &\sim \langle(T_S + \bar{T}_S)\rangle - iC^{\tilde{z}_i\bar{\tilde{z}}_{\bar{j}}}(\mathcal{V})\mathrm{Tr}(\langle\tilde{z}_i\rangle\langle\bar{\tilde{z}}_{\bar{j}}\rangle) \sim \mathcal{O}(1),\\ \langle\sigma_B\rangle &\sim \langle(T_B + \bar{T}_B)\rangle - iC^{\tilde{z}_i\bar{\tilde{z}}_{\bar{j}}}(\mathcal{V})\mathrm{Tr}(\langle\tilde{z}_i\rangle\langle\bar{\tilde{z}}_{\bar{j}}\rangle)\\ &\quad - iC^{a_I\bar{a}_{\bar{J}}}(\mathcal{V})\mathrm{Tr}(\langle a_I\rangle\langle\bar{a}_{\bar{J}}\rangle) \sim e^{f\langle\sigma_S\rangle},\end{aligned} \quad (4)$$

where f is a fraction not too small as compared to 1, and the stabilized values of T_α around the metastable local minimum,

$$\langle\Re e T_S\rangle, \langle\Re e T_B\rangle \sim \mathcal{O}(1). \quad (5)$$

α, a indices correspond to involutively even and odd sectors of $h^{1,1}(CY_3)$ under a holomorphic, isometric involution. If the volume $\mathcal{V}$ of the internal manifold is large in string length units, one sees that one obtains a hierarchy between the stabilized values $\langle\Re e\tau_{S,B}\rangle$ but not $\langle\Re e T_{S,B}\rangle$.

B. Local realization of the model of Sec. II A

We review the local $D3$-$D7$-brane framework presented in [38] which realizes the aforementioned phenomenological supergravity model [(1)–(5)] locally in string theory. In this, we consider type IIB compactified on the orientifold of a Swiss-cheese Calabi-Yau in the large volume scenario (LVS) limit that includes non-(perturbative) α' corrections and nonperturbative instanton corrections in superpotential [41] in addition to a space-time filling $D3$-brane and multiple fluxed stacks of $D7$-branes wrapping the big divisor. We elaborate a little more than what was done in [38] on some algebraic geometric aspects.

The "bottom-up" approach to phenomenological models in the context of D-brane models to realize the SM spectrum was initiated in [42] by considering $D3$-branes on the top of orbifold singularities of $\mathbb{C}^3/\mathbb{Z}_3$ with additional intersecting $D7$-branes (with their world volumes transverse to the respective complex planes). In this model, quarks and one of the Higgs doublets are obtained from strings stretching between different $D3$-branes, while the other Higgs doublet, leptons, and right-handed quark (d_R) are obtained from strings stretching between $D3$- and $D7$-branes; the adjoint gauge fields correspond to open strings starting and ending on the same $D7$-brane. Motivated by this approach, different models were constructed in the context of compact Calabi-Yau compactifications by following configurations of intersecting $D7$-branes wrapping different four cycles (see [43–48] and references therein). With the progress of large volume moduli stabilization [49], realistic constructions reproducing SM spectrum via D-branes were obtained by wrapping $D7$-branes around blown-up cycle(s) [50] (small divisor Σ_s in the geometry of the Swiss-cheese Calabi-Yau orientifold), similar to the techniques used in models of branes at singularities.

The configuration of $D3$-$D7$-branes as described in [38] was also obtained locally in the context of large volume scenarios. However, the setup of [38] is different from the aforementioned large volume scenarios constructs because (i) it considers four stacks of multiple (magnetized) $D7$-branes in groups of three [corresponding to $U(1) \times SU(2)_c$], two [corresponding to $U(1) \times SU(2)_L$], one [corresponding to a $U(1)$], and one [corresponding to another $U(1)$] with the hypercharge corresponding to a linear combination of the four $U(1)$'s wrapping around the big divisor in the rigid limit of the same (given that it was possible to locally stabilize the moduli corresponding to the fluctuations normal to the big divisor Σ_B around which $D7$-branes are wrapped, at null values) but with different choices of two-form fluxes turned on the different two cycles homologously nontrivial from the point of view of this four cycle's homology and not the ambient Swiss-cheese Calabi-Yau. (ii) It takes into account the nonperturbative corrections in the Kähler potential [41] in type IIB Swiss-cheese Calabi-Yau orientifold compactification, not considered in the "large volume scenario" proposed in [49].

Further, similar in spirit to [51–53], by turning on different but small two-form fluxes on the different two cycles homologously nontrivial from the point of view of the big divisor's geometry as a result of which initially adjoint-valued matter fields decompose into bifundamental matter fields corresponding to the SM gauge groups, we provided explicit matrix-valued representations in [38] for $SU(3)_c \times SU(2)_L$ bifundamental first-generation quarks, their right-handed EW-singlet cousins, $SU(2)_L \times U(1)_L$ bifundamental first-generation leptons and Higgs, as well as the right-handed EW-singlet leptonic cousins in [38]. All aforementioned matter fields arise from strings stretched between $D7$-branes stacks with different two-form fluxes turned on. The leptons and quarks get identified with the sermonic superpartners of Wilson line moduli $\mathcal{A}^I$ and the Higgs with the $D3$-brane's position moduli z_i; τ is the axion-dilaton modulus and $\mathcal{G}^a$ are NS-NS and Ramond-Ramond (RR) two-form axions complexified by the axion-dilaton modulus. In the orientifold limit of F theory, one considers an orientifold of the Calabi-Yau involving a holomorphic isometric involution. Though the contribution to the Kähler potential from the matter fields "C_{37}" coming from open strings stretched between the $D3$- and $D7$-branes wrapping Σ_B for Calabi-Yau orientifolds is not known but based on the results for orientifolds of $(T^2)^3$ (see [54]), we guess the following expression: $\frac{|C_{37}|^2}{T_B} \sim \mathcal{V}^{-\frac{1}{18}}|C_{37}|^2$ [using (9)]. Assuming C_{37} to be stabilized at $\mathcal{V}^{-c_{37}}$, $c_{37} > 0$, this contribution would be subdominant relative to other contributions to the Kähler potential. We will, henceforth, ignore the $D3$-$D7$ matter fields.

We will assume that in the coordinate patch (but not globally), $|z_1| \sim \mathcal{V}^{\frac{1}{36}}$, $|z_2| \sim \mathcal{V}^{\frac{1}{36}}$, $|z_3| \sim \mathcal{V}^{\frac{1}{6}}$, the Calabi-Yau looks like the Swiss-cheese $\mathbb{WCP}^4_{1,1,1,6,9}$ [17]. The defining hypersurface for the same is $u_1^{18} + u_2^{18} + u_3^{18} + u_4^3 + u_5^2 - 18\psi \prod_{i=1}^5 u_i - 3\phi(u_1u_2u_3)^6 = 0$. This can be thought of as the following hypersurface in an ambient complex fourfold: $P(x_1,\ldots,x_5;\xi) = 0$ after resolution of the $\mathbb{Z}_3$ singularity [55] (the x_4 and x_5 have been switched relative to [56]; $n = 6$ $\mathbb{CP}^1$ fibration over $\mathbb{CP}^2$ with projective coordinates $x_{1,2,3}$, x_4, x_5 of [56] is equivalent to $n = -6$ with projective coordinates $x_{1,2,3,}$, x_5, x_4; see [57]) with the toric data for the same given by

	x_1	x_2	x_3	x_4	x_5	ξ
Q^1	1	1	1	6	0	9.
Q^2	0	0	0	1	1	2

In the coordinate patch $x_2 \neq 0$ [implying one is away from the $\mathbb{Z}_3$ singular $(0,0,0,x_4,x_5)$ in $\mathbb{WCP}^4_{1,1,1,6,9}$ [17]], $\xi \neq 0$, one sees that the following are the gauge-invariant coordinates: $z_1 = \frac{x_1}{x_2}$, $z_2 = \frac{x_3}{x_2}$, $z_3 = \frac{x_4^2}{x_2^3\xi}$, $z_4 = \frac{x_5^2 x_2^9}{\xi}$. We, henceforth, assume the Calabi-Yau hypersurface to be written in this coordinate patch as $z_1^{18} + z_2^{18} + \mathcal{P}(z_{1,2,3,4};\psi,\phi) = 0$. The divisor $\{x_5 = 0\} \cap \{P(x_{1,2,3,4,5};\xi) = 0\}$ is rigid with $h^{0,0} = 1$ (see [55]) satisfying Witten's unit-arithmetic genus condition and that the Calabi-Yau volume can be written as $\mathrm{vol}(CY_3) = \frac{\tau_4^{3/2}}{18} - \frac{\sqrt{2}\tau_5^{3/2}}{9}$, implying that the "small divisor" Σ_s is $\{x_5 = 0\} \cap \{z_1^{18} + z_2^{18} + \mathcal{P}(z_{1,2,3}, z_4 = 0;\psi,\phi) = 0\}$ and the big divisor Σ_B is $\{x_4 = 0\} \cap \{z_1^{18} + z_2^{18} + \mathcal{P}(z_{1,2,4}, z_3 = 0;\psi,\phi) = 0\}$. Alternatively, using the toric data of [58],

	x_1	x_2	x_3	x_4	x_5	ξ
Q^1	1	1	1	0	0	−3,
Q^2	0	0	0	−2	−3	−1

one can verify that $\{\xi = 0\} \cap \{P'(x_{1,2,3,4,5};\xi) = 0\}$ is the rigid blow-up mode with $h^{0,0} = 1$ (which can be easily verified using COHOMCALG), and one can define gauge-invariant coordinates in the $x_2 \neq 0$, $x_4 \neq 0$ coordinate patch: $z_1 = \frac{x_1}{x_2}$, $z_2 = \frac{x_3}{x_2}$, $z_3 = \frac{(x_5x_1)^2}{x_4^3}$, $z_4 = \frac{(x_6x_1^3)^2}{x_4}$. Interestingly, we found in [38] that the three cycle

$$C_3 : |z_1| \equiv \mathcal{V}^{\frac{1}{36}}, \qquad |z_2| \equiv \mathcal{V}^{\frac{1}{36}}, \qquad |z_3| \equiv \mathcal{V}^{\frac{1}{6}} \tag{6}$$

[the Calabi-Yau can be thought of locally as a complex threefold $\mathcal{M}_3$, which is a T^3 swept out by $\arg z_1$, $\arg z_2$, $\arg z_3$ fibration over a large base $(|z_1|,|z_2|,|z_3|)$]. Precisely apt for application of mirror symmetry as three T dualities *à la* Strominger, Yau, and Zaslow (SYZ), C_3 is almost a special Lagrangian submanifold because it satisfies the requirement that

$$f^*J \approx 0, \qquad \Re e(f^* e^{i\theta}\Omega)|_{\theta=\frac{\pi}{2}} \approx \mathrm{vol}(C_3),$$
$$\Im m(f^* e^{i\theta}\Omega)|_{\theta=\frac{\pi}{2}} \approx 0,$$

where $f: C_3 \to CY_3$. As the defining hypersurface of the Swiss-cheese Calabi-Yau in the $x_2 \neq 0$ coordinate patch will be $z_1^{18} + z_2^{18} + \cdots$, which, near C_3 (implying that the other two coordinates will scale like $\mathcal{V}^{\frac{1}{6}}$, $\mathcal{V}^{\frac{1}{6}} - \mathcal{V}^{\frac{1}{4}}$) receives the most dominant contributions from the monomials z_1^{18} and z_2^{18} it is sufficient to consider $\mathcal{P}_{\Sigma_S}|_{D3|_{\mathrm{near}C_3 \hookrightarrow \Sigma_B}}$, $\mathcal{P}_{\Sigma_B}|_{\mathrm{near}C_3 \hookrightarrow \Sigma_B} \sim z_1^{18} + z_2^{18}$ with the understanding $|\mathcal{P}(z_{1,2,3}, z_4 = 0;\phi,\psi)|_{C_3}, |\mathcal{P}(z_{1,2,4}, z_3;\phi,\psi)|_{C_3} < |z_1^{18} + z_2^{18}|$.

The set of $\mathcal{N} = 1$ chiral coordinates (in particular, the "divisor volume") gets modified in the presence of $D3$- and $D7$-branes [40]. To evaluate the Wilson line moduli contribution in one of the $\mathcal{N} = 1$ chiral coordinates T_B, due to inclusion of four Wilson line moduli on the world volume of space-time filling $D7$-branes wrapped around the big divisor restricted to (nearly) a special Lagrangian submanifold, we constructed distribution harmonic one-forms localized along the mobile space-time filling $D3$-brane (restricted to the three cycle). Here, we review the construction of involutively odd harmonic distribution one-forms in the large volume limit, as given in [38]. [The most nontrivial example of involutions which are meaningful only at large volumes is mirror symmetry implemented as three T dualities in [59] to a Calabi-Yau which locally can be thought of as a T^3 fibration over a (large) base; all Calabi-Yau's with mirrors (in the conventional sense) are expected to have such a local fibration.] Harmonic distribution one-forms can be constructed by integrating $dA_I = (P_{\Sigma_B}(z_{1,2}))^I dz_1 \wedge dz_2$ with ($I = 1, 2, 3, 4$), near $C_3 \hookrightarrow \Sigma_B$; A_I is harmonic only within Σ_B and not at any other generic locus outside Σ_B in the Calabi-Yau manifold. Four such distribution one-forms on Σ_B localized along C_3 corresponding to the location of the $D3$-brane can be written as $A_I \sim \delta(|z_1| - \mathcal{V}^{\frac{1}{36}})\delta(|z_2| - \mathcal{V}^{\frac{1}{36}})[\omega_I(z_1, z_2)dz_1 + \tilde{\omega}_I(z_1, z_2)dz_2]$. Writing $A_I(z_1, z_2) = \omega_I(z_1, z_2)dz_1 + \tilde{\omega}_I(z_1, z_2)dz_2$[1] where $\omega(-z_1, z_2) = \omega(z_1, z_2)$, $\tilde{\omega}(-z_1, z_2) = -\tilde{\omega}(z_1, z_2)$, and $\partial_1\tilde{\omega} = -\partial_2\omega$, one obtains (see [38])

$$\begin{aligned} A_1|_{C_3} &\sim -z_1^{18} z_2^{19} dz_1 + z_1^{19} z_2^{18} dz_2, \\ A_2|_{C_3} &\sim -z_1^{18} z_2 dz_1 + z_2^{18} z_1 dz_2, \\ A_3|_{C_3} &\sim -z_1^{18} z_2^{37} dz_1 - z_2^{18} z_1^{37} dz_1, \\ A_4|_{C_3} &\sim -z_1^{36} z_2^{37} dz_1 + z_2^{36} z_1^{37} dz_2. \end{aligned} \tag{7}$$

1. Yang-Mills coupling constant

We now summarize the discussion on obtaining an $\mathcal{O}(1)$ gauge coupling constant. The Yang-Mills gauge coupling constant squared for the ith gauge group $[i: SU(3), SU(2), U(1)]$ will be given as

$$\frac{1}{g^2_{j=SU(3)\,\text{or}\,SU(2)}} = \Re e(T_{S/B}) + \ln(P(\Sigma_S)|_{D3|_{\Sigma_B}}) + \ln(\bar{P}(\Sigma_S)|_{D3_{\Sigma_B}}) + \mathcal{O}(F_j^2)\tau, \tag{8}$$

where $Re(T_{S/B})$ corresponds to the size of the divisor volume around which $D7$-branes are wrapped, and $F_j^2 = F^\alpha_j F^\beta_j \kappa_{\alpha\beta} + \tilde{F}^\alpha_j \tilde{F}^\beta_j \kappa_{\alpha\beta}$ are the components of the two-form fluxes for the jth stack expanded out in the basis of $i^* w_\alpha$, $w_\alpha \in H^{1,1}_-(CY_3)$, and $\tilde{F}^a_j$ are the components of two-form fluxes for the jth stack expanded out in the basis $\tilde{w}_a \in \text{coker}(H^{(1,2)}_-(CY_3) \to H^{(1,1)}_-(\sigma_B))$. In dilute flux approximation $g_{\rm YM}$ is mainly governed by the size of the divisor volume around which $D7$-branes are wrapped. Using the distribution one-forms of (7), the $\mathcal{N} = 1$ chiral coordinates with the inclusion of mobile $D3$-brane position moduli $z_{1,2}$ (which we identify with the Σ_B coordinates) and mulitple matrix-valued $D7$-branes, Wilson line moduli a_I were guessed in [38]. The quadratic contribution arising in T_B (the big divisor) due to the Wilson line moduli contribution is of the form $i\kappa_4^2\mu_7 C^B_{I\bar{J}} a^I \bar{a}^{\bar{J}}$ with $C^B_{I\bar{J}} = \int_{\Sigma_B} i^*\omega \wedge A^I \wedge \bar{A}^{\bar{J}}$, where $\omega \in H^{(1,1)}_+(\Sigma_B)$. In [38], we estimated the intersection matrices $C^B_{I\bar{J}}$ by constructing harmonic one-forms using Eq. (7). Also, the coefficient of the quadratic term $(\omega_\alpha)_{i\bar{j}} z^i (\bar{z}^{\bar{j}} - \frac{i}{2}(\mathcal{P}_{\tilde{a}})^{\bar{j}}_l \bar{z}^{\tilde{a}} z^l)$ arising in T_B due to inclusion of position moduli z_i was shown in [38] to be $\mathcal{O}(1)$ by calculating $(\omega_B)_{i\bar{j}} \sim (\omega_S)_{i\bar{j}} \sim \mathcal{O}(1)$ near $z_{1,2} \sim \frac{\mathcal{V}^{\frac{1}{36}}}{\sqrt{2}}$ (see [38]). Using the same, it was argued that, in the dilute flux approximation, gauge couplings corresponding to the gauge theories living on stacks of $D7$-branes wrapping the big divisor Σ_B in the large volume limit, will be given by

$$g^{-2}_{\rm YM} \sim \Re e(T_B) \sim \text{vol}(\Sigma_B) + C_{I\bar{J}} a_I \bar{a}_{\bar{J}} + \text{H.c.} \sim \mathcal{V}^{\frac{1}{18}} \sim O(1)$$

(justified by the partial cancellation between Σ_B and $C_{I\bar{J}} a_I \bar{a}_{\bar{J}}$ with some fine-tuning).

2. Stabilized potential of $\mathcal{N} = 1$ local large volume D3-D7 setup

As we do not have a global picture, we are ourselves with a local bulk and open-string moduli stabilization near (6). We showed in [38] that near (6), the moduli can be stabilized as under

$$\begin{aligned} &\text{vol}(\Sigma_S) \sim \mathcal{V}^{\frac{1}{18}}, \qquad \text{vol}(\Sigma_B) \sim \mathcal{V}^{\frac{2}{3}}; \\ &\mathcal{G}^a \sim \frac{\pi}{\mathcal{O}(1)k^a(\sim\mathcal{O}(10))} M_P; \\ &|z_{1,2}| \equiv \mathcal{V}^{\frac{1}{36}} M_P, \qquad |z_3| \equiv \mathcal{V}^{\frac{1}{6}} M_P; \\ &|a_1| \equiv \mathcal{V}^{-\frac{2}{9}} M_P, \qquad |a_2| \equiv \mathcal{V}^{-\frac{1}{3}} M_P, \\ &|a_3| \equiv \mathcal{V}^{-\frac{13}{18}} M_P, \qquad |a_4| \equiv \mathcal{V}^{-\frac{11}{9}} M_P; \\ &\zeta^{A=1,\ldots,h^{0,2}_-(\Sigma_B|_{C_3})} \equiv 0\,(\text{implying rigidity of the nonrigid}\,\Sigma_B), \end{aligned} \tag{9}$$

such that $\Re e T_S \sim \Re e T_B \sim \mathcal{V}^{\frac{1}{18}}$ and implying the possibility of obtaining a local metastable de Sitter–like minimum corresponding to the positive minimum of the potential $e^K G^{T_S \bar{T}_S} |D_{T_S} W|^2$ near (9), and realizing (5) and thereby the supergravity model of Sec. II A for $\mathcal{V} \sim 10^5$ in $l_s = 1$ units.

The Kähler potential relevant to all the calculations (using modified $\mathcal{N} = 1$ chiral coordinates) in this paper [without being careful about $\mathcal{O}(1)$ constant factors] is given as under [38],

[1]Intuitively, these distribution one-forms could be thought of as the holomorphic square root of a Poincaré dual of a four cycle.

$$K \sim -2\ln\left(a_B\left[\frac{T_B+\bar{T}_B}{M_P} - \mu_3(2\pi\alpha')^2\frac{\{|z_1|^2+|z_2|^2+z_1\bar{z}_2+z_2\bar{z}_1\}}{M_P^2} + \mathcal{V}^{\frac{10}{9}}\frac{|a_1|^2}{M_P^2} + \mathcal{V}^{\frac{11}{18}}\frac{(a_1\bar{a}_2+\text{H.c.})}{M_P^2} + \mathcal{V}^{\frac{1}{9}}\frac{|a_2|^2}{M_P^2} + \mathcal{V}^{\frac{29}{18}}\frac{(a_1\bar{a}_3+\text{H.c.})}{M_P^2} + \mathcal{V}^{\frac{10}{9}}\frac{(a_2\bar{a}_3+\text{H.c.})}{M_P^2} + \mathcal{V}^{\frac{19}{9}}\frac{|a_3|^2}{M_P^2} + \mathcal{V}^{\frac{19}{9}}\frac{(a_1\bar{a}_4+a_4\bar{a}_1)}{M_P^2} + \mathcal{V}^{\frac{29}{18}}\frac{(a_2\bar{a}_4+a_4\bar{a}_2)}{M_P^2} + \mathcal{V}^{\frac{47}{18}}\frac{(a_3\bar{a}_4+a_4\bar{a}_3)}{M_P^2} + \mathcal{V}^{\frac{28}{9}}\frac{|a_4|^2)}{M_P^2}\right]^{3/2} - a_S\left(\frac{T_S+\bar{T}_S}{M_P} - \mu_3\left(2\pi\alpha'\right)^2\frac{\{|z_1|^2+|z_2|^2+z_1\bar{z}_2+z_2\bar{z}_1\}}{M_P^2}\right)^{3/2} + \sum n^0_\beta(\ldots)\right), \tag{10}$$

and the ED3/ED1-generated nonperturbative superpotential used in [38] is given by

$$W \sim (\mathcal{P}_{\Sigma_S}|_{D3|_{\text{near}\,C_3\hookrightarrow\Sigma_B}} \sim z_1^{18}+z_2^{18})^{n^s}\sum_{m_a} e^{i\tau\frac{m^2}{2}}e^{in^sG^am_a}e^{in^sT_s}, \tag{11}$$

which is like (3) assuming G^a, τ has been stabilized. The genus-zero Gopakumar-Vafa invariants (which, for projective varieties, are very large) prefix the $h^{1,1}_-$-valued real axions b^a, c^a. In general, there are no known globally defined involutions valid for all Calabi-Yau volumes, for which $h^{1,1}_-(CY_3)\neq 0$, $h^{0,1}_-(\Sigma_B)\neq 0$. However, as mentioned earlier, in the spirit of the involutive mirror symmetry implemented *à la* the SYZ prescription in terms of a triple of T dualities along a local T^3 in the large volume limit, we argued in [60], e.g., $z_1\to -z_1$ would, restricted to C_3, generate nonzero

$$h^{1,1}_-\begin{pmatrix} T^3(\arg z_{1,2,3}) \to \mathcal{M}_3(z_{1,2,3}) \\ \downarrow \\ M_3(|z_1|,|z_2|,|z_3|)\end{pmatrix}.$$

An example of holomorphic involutions near C_3 not requiring a large Calabi-Yau volume has been discussed in [38]. However, even if $h^{1,1}_-=0$, one can self-consistently stabilize c^a, b^a to zero and σ_s, σ_b to $\mathcal{V}^{\frac{1}{18}}$, $\mathcal{V}^{\frac{2}{3}}$ such that the Kähler potential continues to be stabilized at $-2\ln\mathcal{V}$.

The evaluation of "physical"/normalized Yukawa couplings, soft SUSY breaking parameters, and various three-point vertices needs the matrix generated from the mixed double derivative of the Kähler potential to be a diagonalized matrix. After diagonalization, the corresponding eigenvectors of the same were found in [38] to be given by

$$\begin{aligned}
&\mathcal{A}_4 \sim a_4+\mathcal{V}^{-\frac{3}{5}}a_3+\mathcal{V}^{-\frac{6}{5}}a_1+\mathcal{V}^{-\frac{9}{5}}a_2+\mathcal{V}^{-2}(z_1+z_2),\\
&\mathcal{A}_3 \sim -a_3+\mathcal{V}^{-\frac{3}{5}}a_4-\mathcal{V}^{-\frac{3}{5}}a_1-\mathcal{V}^{-\frac{7}{5}}a_2+\mathcal{V}^{-\frac{8}{5}}(z_1+z_2),\\
&\mathcal{A}_1 \sim a_1-\mathcal{V}^{-\frac{3}{5}}a_3+\mathcal{V}^{-1}a_2-\mathcal{V}^{-\frac{6}{5}}a_4+\mathcal{V}^{-\frac{6}{5}}(z_1+z_2),\\
&\mathcal{A}_2 \sim -a_2-\mathcal{V}^{-1}a_1+\mathcal{V}^{-\frac{7}{5}}a_3-\mathcal{V}^{-\frac{3}{5}}(z_1+z_2),\\
&\mathcal{Z}_2 \sim -\frac{(z_1+z_2)}{\sqrt{2}}-\mathcal{V}^{-\frac{6}{5}}a_1+\mathcal{V}^{-\frac{3}{5}}a_2+\mathcal{V}^{-\frac{8}{5}}a_3+\mathcal{V}^{-2}a_4,\\
&\mathcal{Z}_1 \sim \frac{(z_1-z_2)}{\sqrt{2}}-\mathcal{V}^{-\frac{6}{5}}a_1+\mathcal{V}^{-\frac{3}{5}}a_2+\mathcal{V}^{-\frac{8}{5}}a_3+\mathcal{V}^{-2}a_4.
\end{aligned}$$

For $\mathcal{V}=10^5$, the numerical eigenvalues are estimated to be

$$\begin{aligned}
&K_{\mathcal{Z}_1\mathcal{Z}_1}\sim 10^{-5}, \quad K_{\mathcal{Z}_2\mathcal{Z}_2}\sim 10^{-3}, \quad K_{\mathcal{A}_1\mathcal{A}_1}\sim 10^{4},\\
&K_{\mathcal{A}_2\mathcal{A}_2}\sim 10^{-2}, \quad K_{\mathcal{A}_3\mathcal{A}_3}\sim 10^{7}, \quad K_{\mathcal{A}_4\mathcal{A}_4}\sim 10^{12}.
\end{aligned} \tag{12}$$

3. Mass scales of SM-like particles

The effective Yukawa couplings can be calculated using $\hat{Y}^{\text{eff}}_{C_iC_jC_k}\equiv\frac{e^{\frac{K}{2}}Y^{\text{eff}}_{C_iC_jC_k}}{\sqrt{K_{C_i\bar{C}_i}K_{C_j\bar{C}_j}K_{C_k\bar{C}_k}}}$, C_i being an open-string modulus, which for us is $\delta\mathcal{Z}_{1,2}$, $\delta\mathcal{A}_{1,2,3,4}$, where $Y^{\text{eff}}_{\mathcal{Z}_i\mathcal{A}_I\mathcal{A}_J}$ is given by the $\mathcal{O}(\mathcal{Z}_i)$ coefficient in the mass term $e^{\frac{K}{2}}\mathcal{D}_{\bar{\mathcal{A}}_I}D_{\bar{\mathcal{A}}_J}\bar{W}\bar{\chi}^{\mathcal{A}_I}\chi^{\mathcal{A}_J}$ in the $\mathcal{N}=1$ SUGRA action of [39]. By estimating in the large volume limit, all possible Yukawa couplings corresponding to four Wilson line moduli and showing that the renormalization group (RG) flow of the effective physical Yukawa's change almost by $\mathcal{O}(1)$ under a RG flow from the string scale down to the EW scale [38], we see that for $\mathcal{V}\sim 10^5$, $\langle\mathcal{Z}_i\rangle\sim 246$ GeV:

$$\frac{\mathcal{O}(\mathcal{Z}_i)\text{term in }e^{\frac{K}{2}}\mathcal{D}_{\mathcal{A}_1}D_{\mathcal{A}_3}W}{\sqrt{K_{\mathcal{Z}_i\bar{\mathcal{Z}}_i}K_{\mathcal{A}_1\bar{\mathcal{A}}_1}K_{\mathcal{A}_3\bar{\mathcal{A}}_3}}}\equiv\hat{Y}^{\text{eff}}_{\mathcal{Z}_i\mathcal{A}_1\mathcal{A}_3}\sim 10^{-3}\times\mathcal{V}^{-\frac{4}{9}},$$

giving $\langle\mathcal{Z}_i\rangle\hat{Y}_{\mathcal{Z}_1\mathcal{A}_1\mathcal{A}_3}\sim$ MeV—about the mass of the electron

$$\frac{\mathcal{O}(\mathcal{Z}_i)\text{term in }e^{\frac{K}{2}}\mathcal{D}_{\mathcal{A}_2}D_{\mathcal{A}_4}W}{\sqrt{K_{\mathcal{Z}_i\bar{\mathcal{Z}}_i}K_{\mathcal{A}_2\bar{\mathcal{A}}_2}K_{\mathcal{A}_4\bar{\mathcal{A}}_4}}}\equiv\hat{Y}^{\text{eff}}_{\mathcal{Z}_i\mathcal{A}_2\mathcal{A}_4}\sim 10^{-\frac{5}{2}}\times\mathcal{V}^{-\frac{4}{9}}$$

giving $\langle z_i\rangle \hat{Y}_{\mathcal{Z}_i\mathcal{A}_2\mathcal{A}_4} \sim 10$ MeV—close to the mass of the up quark. The above shows that fermionic superpartners of $\mathcal{A}_1$ and $\mathcal{A}_3$ correspond, respectively, to the first generation of left-handed $SU(2)$ and right-handed $U(1)$ leptons, while fermionic superpartners of $\mathcal{A}_2$ and $\mathcal{A}_4$ correspond, respectively, to left-handed $SU(2)$ and right-handed $U(1)$ quarks. The diagonalized basis (12) was shown to also work out for appropriately chosen matrix-valued a_I and z_i for multiple fluxed $D7$-brane stacks.

4. Computation of soft terms

By using the appropriate $\mathcal{N}=1$ coordinates as obtained in [40] due to the presence of a single $D3$-brane and a single $D7$-brane wrapping, the four cycle (big divisor Σ_B in a Swiss-cheese Calabi-Yau) along with $D7$-brane fluxes, the soft SUSY breaking parameters were calculated in [38]. The value of scalar masses identified with the masses of squarks and leptons, so obtained, turns out to be quite high, but at the same time, one gets one light Higgs, thus, indicating the possibility of a "split SUSY-like scenario" in a local large volume $D3/D7$ model.

We briefly review the evaluation of various soft supersymmetric as well as supersymmetry breaking parameters in the model involving four Wilson line moduli as described in [38]. The various soft terms are calculated by power series expansion of the superpotential as well as the Kähler potential,

$$W=\hat{W}(\Phi)+\mu(\Phi)\mathcal{Z}_I\mathcal{Z}_J+\frac{1}{6}Y_{IJK}(\Phi)\mathcal{M}^I\mathcal{M}^J\mathcal{M}^K+\cdots,$$
$$K=\hat{K}(\Phi,\bar{\Phi})+K_{I\bar{J}}(\Phi,\bar{\Phi})\mathcal{M}^I\mathcal{M}^{\bar{J}}+Z(\Phi,\bar{\Phi})\mathcal{M}^I\mathcal{M}^{\bar{J}}+\cdots, \tag{13}$$

where $\mathcal{M}^I=(\mathcal{Z}^I,\mathcal{A}^I)$. The soft SUSY breaking parameters are calculated by expanding the $\mathcal{N}=1$ supergravity potential, $V=e^K(K^{I\bar{J}}D_IWD_{\bar{J}}\bar{W}-3|W|^2)$ in the powers of matter fields $\mathcal{M}^I$ after expanding the superpotential and Kähler potential according to Eq. (13). In gravity-mediated supersymmetry breaking, SUSY gets spontaneously broken in the bulk sector by giving a vacuum expectation value to auxiliary F terms. Hence, to begin with, one needs to evaluate the bulk F terms, which, in turn, entails evaluating the bulk metric. Writing the Kähler sector of the Kähler potential in terms of the bulk moduli as

$$K\sim-2\ln\Bigg[(\sigma_B+\bar{\sigma}_B-\gamma K_{\rm geom})^{\frac{3}{2}}-(\sigma_S+\bar{\sigma}_S-\gamma K_{rm{\rm geom}})^{\frac{3}{2}}+\sum_{\beta\in H_2^-(CY_3)}n_\beta^0\sum_{(n,m)}\cos\left(ink\cdot(G-\bar{G})g_s+mk\cdot(G+\bar{G})\right)\Bigg], \tag{14}$$

disregarding $K_{\rm geom}$ (introduced due to the presence of a mobile space-time filling $D3$-brane) in the large volume limit (see [61,62]) and working near $\sin\left(ink\cdot(G-\bar{G})g_s+mk\cdot(G+\bar{G})\right)=0$ corresponding to a local minimum—using the stabilized vacuum expectation value (VEV) of $\sigma_{S/B}$ and $\mathcal{G}^{S,B}$ as given above Eq. (10)—generated the following components of the bulk metric's inverse in [38]:

$$G^{m\bar{n}}\sim\begin{pmatrix}\mathcal{V}^{\frac{37}{36}} & \mathcal{V}^{\frac{13}{18}} & 0 & 0\\ \mathcal{V}^{\frac{13}{18}} & \mathcal{V}^{\frac{4}{3}} & 0 & 0\\ 0 & 0 & \mathcal{O}(1) & \mathcal{O}(1)\\ 0 & 0 & \mathcal{O}(1) & \mathcal{O}(1)\end{pmatrix}. \tag{15}$$

Given that bulk F terms are defined as [51], $F^m=e^{\frac{K}{2}}G^{m\bar{n}}D_{\bar{n}}\bar{W}$, one obtained in [38]:

$$F^{\sigma_S}\sim\mathcal{V}^{-\frac{n^s}{2}+\frac{1}{36}}M_P^2,\quad F^{\sigma_B}\sim\mathcal{V}^{-\frac{n^s}{2}-\frac{5}{18}}M_P^2,\quad F^{G^a}\sim\mathcal{V}^{-\frac{n^s}{2}-1}M_P^2. \tag{16}$$

Hence, after spontaneous supersymmetry breaking in the bulk, the gravitino mass is given by

$$m_{3/2}=e^K|W|^2\sim\mathcal{V}^{-\frac{n^s}{2}-1}M_P. \tag{17}$$

The gaugino mass is given as

$$m_{\tilde{g}}=\frac{F^m\partial_mT_B}{ReT_B}\lesssim\mathcal{V}^{\frac{2}{3}}m_{3/2}. \tag{18}$$

The analytic form of the scalar masses obtained via spontaneous symmetry breaking is given as [51] $m_I^2=(m_{\frac{3}{2}}^2+V_0)-F^{\bar{m}}F^n\partial_{\bar{m}}\partial_n\log K_{I\bar{I}}$. These were calculated in [38] to yield

$$m_{\mathcal{Z}_i}\sim\mathcal{V}^{\frac{59}{72}}m_{3/2},\qquad m_{\mathcal{A}_1}\sim\sqrt{\mathcal{V}}m_{3/2}, \tag{19}$$

implying a nonuniversality in the open-string moduli masses. Further, in [38] we showed the universality in the trilinear A couplings [51],

$$A_{\mathcal{I}\mathcal{J}\mathcal{K}}=F^m(\partial_mK+\partial_m\ln Y_{\mathcal{I}\mathcal{J}\mathcal{K}}+\partial_m\ln(K_{\mathcal{I}\bar{\mathcal{I}}}K_{\mathcal{J}\bar{\mathcal{J}}}K_{\mathcal{K}\bar{\mathcal{K}}}))$$
$$\sim\mathcal{V}^{\frac{37}{36}}m_{3/2}\sim\hat{\mu}_{\mathcal{Z}_1\mathcal{Z}_2}. \tag{20}$$

The physical Higgsino mass parameter $\hat{\mu}_{\mathcal{Z}_1\mathcal{Z}_2}$ turned out to be given by

$$\hat{\mu}_{\mathcal{Z}_1\mathcal{Z}_2}=\frac{e^{\frac{K}{2}}\mu_{\mathcal{Z}_1\mathcal{Z}_2}}{\sqrt{K_{\mathcal{Z}_1\bar{\mathcal{Z}}_1}K_{\mathcal{Z}_2\bar{\mathcal{Z}}_2}}}\sim\mathcal{V}^{\frac{19}{18}}m_{3/2}. \tag{21}$$

Further,

$$(\hat{\mu}B)_{\mathcal{Z}_1\mathcal{Z}_2} = \frac{e^{-i\arg(W)+\frac{K}{2}}}{\sqrt{K_{\mathcal{Z}_1\bar{\mathcal{Z}}_1}K_{\mathcal{Z}_2\bar{\mathcal{Z}}_2}}}F^m(\partial_m K\mu_{\mathcal{Z}_1\mathcal{Z}_2} + \partial_m\mu_{\mathcal{Z}_1\mathcal{Z}_2} - \mu_{\mathcal{Z}_1\mathcal{Z}_2}\partial_m \ln(K_{\mathcal{Z}_1\bar{\mathcal{Z}}_1}K_{\mathcal{Z}_2\bar{\mathcal{Z}}_2})) \sim \hat{\mu}_{\mathcal{Z}_1\mathcal{Z}_2}(F^m\partial_m K + F^{\sigma_s} - F^m\partial_m \ln(K_{\mathcal{Z}_1\bar{\mathcal{Z}}_1}K_{\mathcal{Z}_2\bar{\mathcal{Z}}_2})) \sim \mathcal{V}^{\frac{19}{18}+\frac{37}{36}}m_{3/2}^2 \sim \hat{\mu}^2_{\mathcal{Z}_1\mathcal{Z}_2}, \tag{22}$$

an observation which will be very useful in obtaining a light Higgs of mass 125 GeV.

5. Realizing a light SM-like Higgs

We calculated in [38] the mass of a light Higgs formed by the linear combination of two Higgs doublets (using the prescription as given in [1] to realize split SUSY) by first calculating the masses of the latter, which, after soft supersymmetry breaking, are given by $M_{H_{u,d}} = (m^2_{\mathcal{Z}_{1,2}} + \hat{\mu}^2_{\mathcal{Z}_1\mathcal{Z}_2})^{1/2}$, and, thereafter, using the RG solution to the Higgs mass discussed in [37], we obtained the contribution of Higgs doublets as well as the Higgsino mass parameter $\hat{\mu}_{\mathcal{Z}_1\mathcal{Z}_2}$ at the EW scale. The Higgs mass eigenstates are defined as

$$H_1 = D_{h_{11}}H_u + D_{h_{12}}H_d, \qquad H_2 = D_{h_{21}}H_u + D_{h_{22}}H_d, \tag{23}$$

where

$$D_h = \begin{pmatrix} \cos\frac{\theta_h}{2} & -\sin\frac{\theta_h}{2}e^{-i\phi_h} \\ \sin\frac{\theta_h}{2}e^{i\phi_h} & \cos\frac{\theta_h}{2} \end{pmatrix},$$

$D_h^\dagger M_h^2 D_h = \mathrm{diag}(M^2_{H_1}, M^2_{H_2})$, and $\tan\theta_h = \frac{2|M^2_{h_{21}}|}{M^2_{h_{11}} - M^2_{h_{22}}}$ for a particular range of $\frac{-\pi}{2} \leq \theta_h \leq \frac{\pi}{2}$.

The RG solution to the Higgs mass formed after soft supersymmetry breaking in the large-tan β (but less than 50) limit are given [37,38] [assuming that $m^2_{\mathcal{Z}_2}(M_s) \equiv m_0^2 \equiv \mathcal{V}^{\frac{59}{72}}m_{\frac{3}{2}}$, implying $\delta_2 = 0$ but $\delta_{1,3,4} \neq 0$, and nonuniversality with respect to both $D3$-brane position moduli masses ($m_{Z_{1,2}}$) given by δ_1] as

$$\hat{\mu}^2(\mathrm{EW}) \equiv -\left[-m_0^2 - (0.01)(n^s)^2\hat{\mu}^2_{\mathcal{Z}_1\mathcal{Z}_2} + (0.32)\mathcal{V}^{\frac{4}{3}}m^2_{3/2} - 1/2M^2_{\mathrm{EW}} + (0.03)\mathcal{V}^{\frac{2}{3}}n^s\hat{\mu}_{\mathcal{Z}_1\mathcal{Z}_2}m_{3/2} + \frac{19\pi}{2200}S_0\right],$$

$$m^2_{H_u}(\mathrm{EW}) \equiv m_0^2(1+\delta_1) + \frac{1}{2}M^2_{\mathrm{EW}} + m_0^2 - (0.03)\mathcal{V}^{\frac{2}{3}}n^s\hat{\mu}_{\mathcal{Z}_1\mathcal{Z}_2}m_{3/2} + (0.01)(n^s)^2\hat{\mu}^2_{\mathcal{Z}_1\mathcal{Z}_2},$$

$$m^2_{H_d}(\mathrm{EW}) \equiv 2m_0^2 - (0.06)\mathcal{V}^{\frac{2}{3}}n^s\hat{\mu}_{\mathcal{Z}_1\mathcal{Z}_2}m_{3/2} + \frac{1}{2}M^2_{\mathrm{EW}} - \frac{19\pi}{1100}S_0,$$

where S_0 is a hypercharge weighted sum of the squared soft scalar mass having value around m_0^2. Assuming $\hat{\mu}B \equiv \xi\hat{\mu}_{Z_iZ_j}(\xi \equiv O(1))$, the Higgs mass matrix is given as

$$\begin{pmatrix} m^2_{H_u} & \hat{\mu}B \\ \hat{\mu}B & m^2_{H_d} \end{pmatrix} \sim \begin{pmatrix} m^2_{H_u} & \xi\hat{\mu}^2 \\ \xi\hat{\mu}^2 & m^2_{H_d} \end{pmatrix},$$

and the eigenvalues are given by $\frac{1}{2}(m^2_{H_u} + m^2_{H_d} \pm \sqrt{(m^2_{H_u} - m^2_{H_d})^2 + 4\xi^2\hat{\mu}^4})$. Using Eq. (24), for $\delta_1 = \mathcal{O}(0.1)$ and $\mathcal{O}(1)\,n^s$, we have

$$m^2_{H_u} + m^2_{H_d} \sim m_0^2 - 0.06S_0 + \cdots,$$
$$m^2_{H_u} - m^2_{H_d} \sim m_0^2 + 0.06S_0 + \cdots,$$
$$\hat{\mu}^2_{\mathcal{H}_u\mathcal{H}_d} \sim m_0^2 - 0.03S_0 + \cdots$$

Utilizing the above, one sees that the eigenvalues are

$$m^2_{H_{1,2}} = m_0^2 - 0.06S_0 + \cdots \pm \sqrt{(m_0^2 + 0.06S_0 + \cdots)^2 + 4\xi^2(m_0^2 - 0.03S_0)^2}.$$

Considering $S_0 \sim -4.2m_0^2$ and $\xi^2 \sim \frac{1}{5} + \frac{1}{16}\frac{m^2_{\mathrm{EW}}}{m_0^2}$, we obtain one light Higgs (corresponding to the negative sign of the square root) of order 125 GeV and one heavy Higgs (corresponding to the positive sign of the square root), whereas the squared Higgsino mass parameter $\hat{\mu}_{\mathcal{Z}_1\mathcal{Z}_2}$ then turns out to be heavy with a value at the EW scale of around $\mathcal{V}m_{3/2}$.

6. Realization of a μ-split-like SUSY

We summarized above the different mass scales corresponding to different supersymmetric particles as mentioned in the above paragraphs and actually calculated in [38] by considering Calabi-Yau volume $\mathcal{V} = 10^5$ [the justification behind constraining a value of Calabi-Yau $\mathcal{V}$ to be $\mathcal{O}(10^5)$ was based on the right identification of Wilson line moduli and position moduli with a SM particle spectrum]. The gravitino appears to be the lightest supersymmetric particle with mass around 10^8 GeV. The sfermion masses corresponding to the first generation of quarks and leptons (identifiable as Wilson line moduli mass in our framework as mentioned above) are very heavy, of the order 10^{10} GeV at the string scale. Similarly, the gaugino masses also turn out to be heavy, of the order 10^{11} GeV. However, the Higgsino masses are heavier, of the order 10^{13} GeV. One of the Higgs doublets was shown to have mass of the order 125 GeV, thus, showing the possibility of realizing a μ-split-like SUSY scenario (though there is a "split" between the mass of a Higgsino, and the gaugino and sfermions at very high energy scale, the SM fermions are light) in the context of our local LVS $D3$-$D7$.

TABLE I. Mass scales of the first generation of the SM as well as supersymmetric and soft SUSY breaking parameters.

Quark mass	$M_q \sim O(10)$ MeV
Lepton mass	$M_l \sim \mathcal{O}(1)$ MeV
Gravitino mass	$m_{\frac{3}{2}} \sim \mathcal{V}^{-\frac{n^s}{2}-1} M_P;\ n_s = 2$
Gaugino mass	$M_{\tilde{g}} \sim \mathcal{V}^{\frac{2}{3}} m_{\frac{3}{2}}$
(Lightest) neutralino/chargino mass	$M_{\chi_3^0/\chi_3^{\pm}} \sim \mathcal{V}^{\frac{2}{3}} m_{\frac{3}{2}}$
$D3$-brane position moduli (Higgs) mass	$m_{\mathcal{Z}_i} \sim \mathcal{V}^{\frac{59}{72}} m_{\frac{3}{2}}$
Wilson line moduli (sfermion mass)	$m_{\tilde{\mathcal{A}}_I} \sim \mathcal{V}^{\frac{1}{2}} m_{\frac{3}{2}}$ $I = 1, 2, 3, 4$
A terms	$A_{pqr} \sim n^s \mathcal{V}^{\frac{37}{36}} m_{\frac{3}{2}}$ $\{p, q, r\} \in \{\tilde{\mathcal{A}}_I, \mathcal{Z}_i\}$
Physical μ terms (Higgsino mass)	$\hat{\mu}_{\mathcal{Z}_i \mathcal{Z}_j} \sim \mathcal{V}^{\frac{37}{36}} m_{\frac{3}{2}}$
Physical $\hat{\mu}B$ terms	$(\hat{\mu}B)_{\mathcal{Z}_1 \mathcal{Z}_2} \sim \mathcal{V}^{\frac{37}{18}} m_{\frac{3}{2}}^2$

The fine-tuning involved in the hypercharge weighted sum of soft scalar masses (S_0) as well as the $\mathcal{O}(1)$ proportionality constant between the Higgsino mass parameter squared μ^2 and the soft SUSY parameter μB to obtain a Higgs of the order 125 GeV seems acceptable at such high energy scales. The results of mass scales of all SM as well as superpartners are summarized in Table I also.

7. Modified $\mathcal{N} = 1$ gauged supergravity action in the case of multiple D7-branes

We will be using the following terms (written out in four-component notation or their two-component analogs and utilizing/generalizing the results of [40]) in the $\mathcal{N} = 1$ gauged supergravity action of Wess and Bagger [39] with the understanding that $m_{\text{moduli/modulini}} \ll m_{\text{KK}} (\sim \frac{M_s}{\mathcal{V}^{\frac{1}{6}}}|_{\mathcal{V}\sim 10^{5/6}} \sim 10^{14}\ \text{GeV}),\ M_s = \frac{M_P}{\sqrt{\mathcal{V}}}|_{\mathcal{V}\sim 10^{5/6}} \sim 10^{15}\ \text{GeV}$ and that for multiple $D7$-branes, the non-Abelian gauged isometry group[2] corresponding to the Killing vector $6i\kappa_4^2 \mu_7 (2\pi\alpha') Q_B \partial_{T_B},\ Q_B = (2\pi\alpha') \int_{\Sigma_B} i^* \omega_B \wedge P_- \tilde{f}$ arising due to the elimination of the two-form axions $D_B^{(2)}$ in favor of the zero-form axions ρ_B under the Kaluza-Klein (KK) reduction of the ten-dimensional four-form axion [40] [which results in a modification of the covariant derivative of T_B by an additive shift given by $6i\kappa_4^2 \mu_7 (2\pi\alpha') \text{Tr}(Q_B A_\mu)$] can be identified with the SM group (i.e., A_μ is the SM-like adjoint-valued gauge field [39]):

$$
\begin{aligned}
\mathcal{L} = {} & g_{YM} g_{T_B \bar{\mathcal{J}}} \text{Tr}(X^{T_B} \bar{\chi}_L^{\bar{\mathcal{J}}} \lambda_{\tilde{g},R}) + i g_{\mathcal{I}\bar{\mathcal{J}}} \text{Tr}\left(\bar{\chi}_L^{\bar{\mathcal{I}}} \left[\partial\!\!\!/ \chi_L^{\mathcal{I}} + \Gamma^i_{Mj} \partial\!\!\!/ a^M \chi_L^{\mathcal{J}} + \frac{1}{4}(\partial_{a_M} K \partial\!\!\!/ a_M - \text{c.c.}) \chi_L^{\mathcal{I}}\right]\right) \\
& + \frac{e^{\frac{K}{2}}}{2} (\mathcal{D}_{\bar{\mathcal{I}}} D_{\mathcal{J}} \bar{W}) \text{Tr}(\chi_L^{\mathcal{I}} \chi_R^{\mathcal{J}}) + g_{T_B \bar{T}_B} \text{Tr}[(\partial_\mu T_B - A_\mu X^{T_B})(\partial^\mu T_B - A^\mu X^{T_B})^\dagger] \\
& + g_{T_B \mathcal{J}} \text{Tr}(X^{T_B} A_\mu \bar{\chi}_L^{\mathcal{J}} \gamma^\nu \gamma^\mu \psi_{\nu,R}) + \bar{\psi}_{L,\mu} \sigma^{\rho\lambda} \gamma^\mu \lambda_{\tilde{g},L} F_{\rho\lambda} + \bar{\psi}_{L,\mu} \sigma^{\rho\lambda} \gamma^\mu \lambda_{\tilde{g},L} W_\rho^+ W_\lambda^- \\
& + \text{Tr}\left[\bar{\lambda}_{\tilde{g},L} A\!\!\!/ \left(6\kappa_4^2 \mu_7 (2\pi\alpha') Q_B K + \frac{12\kappa_4^2 \mu_7 (2\pi\alpha') Q_B v^B}{\mathcal{V}}\right) \lambda_{\tilde{g},L}\right] \\
& + \frac{e^K G^{T_B \bar{T}_B}}{\kappa_4^2} 6i\kappa_4^2 (2\pi\alpha') \text{Tr}[Q_B A^\mu \partial_\mu (\kappa_4^2 \mu_7 (2\pi\alpha')^2 C^{I\bar{J}} a_I \bar{a}_{\bar{J}})] + \text{H.c.} \\
& - \frac{f_{ab}}{4} F^a_{\mu\nu} F^{bmu\nu} + \frac{1}{8} f_{ab} \epsilon^{\mu\nu\rho\lambda} F^a_{\mu\nu} F^b_{\rho\lambda} - \frac{i\sqrt{2}}{4} g \partial_{i/I} f_{ab} \text{Tr}\left(\frac{12\kappa_4^2 \mu_7 (2\pi\alpha') Q_B^a v^B}{\mathcal{V}} \bar{\lambda}^b_{\tilde{g},L} \chi_R^{i/I}\right) + \text{H.c.} \\
& - \frac{\sqrt{2}}{4} \partial_{i/I} f_{ab} \text{Tr}(\bar{\lambda}^a_{\tilde{g},R} \sigma^{\mu\nu} \chi_L^{i/I}) F^b_{\mu\nu} + \text{H.c.}
\end{aligned}
\tag{24}
$$

III. *CP*-VIOLATING PHASES

In this section, we explain the possible origin of CP-violating phases in the $\mathcal{N} = 1$ gauged supergravity limit of the large volume $D3/D7$ μ-split SUSY model. The electric dipole moment of a spin-$\frac{1}{2}$ particle is defined by the effective CP-violating dimension-five operator given as $\mathcal{L}_I = -\frac{i}{2} d_f \bar{\psi} \sigma_{\mu\nu} \gamma_5 \psi F^{\mu\nu}$. Given that the effective operator is nonrenormalizable, the same can be realized at the loop level provided the theory contains a source of CP violation. In the Standard Model, CP-violating phases, in general,

[2] As explained in [40], one of the two Pecci-Quinn shift symmetries along the RR two-form axions c^a and the four-form axion ρ_B gets gauged due to the dualization of the Green-Schwarz term $\int_{\mathbf{R}^{1,3}} dD_B^{(2)} \wedge A$ coming from the KK reduction of the Chern-Simons term on $\Sigma_B \cup \sigma(\Sigma_B) - D_B^{(2)}$ being an RR two-form axion. In the presence of fluxes for multiple $D7$-brane fluxes, the aforementioned Green-Schwarz is expected to be modified to $\text{Tr}(Q_B \int_{\mathbf{R}^{1,3}} dD_B^{(2)} \wedge A)$, which, in turn, after dualization modifies the covariant derivative of T_B and, hence, the Killing isometry.

appear from the CKM phases in the quark mass matrix but the same get a nonzero contribution only at three-loop level in the Standard Model. However, in supersymmetric theories, instead of the CKM phase generated in the Standard Model, one can consider the new phases appearing from complex parameters of soft SUSY breaking terms, complex effective Yukawa couplings, as well as supersymmetric mass terms.

We consider the existence of nonzero phases appearing from complex effective Yukawa couplings present in the $\mathcal{N}=1$ gauged supergravity action. As discussed in [38], the position as well as Wilson line moduli identification with SM-like particles generate effective Yukawa couplings including R-parity conserving as well as R-parity violating ones in the context of the $\mathcal{N}=1$ gauged supergravity action [38], and the solution of RG evolution of effective Yukawa couplings at one-loop level yields

$$\hat{Y}_{\Lambda\Sigma\Delta}(t) \sim \hat{Y}_{\Lambda\Sigma\Delta}(M_s) \prod_{(a)=1}^{3} (1+\beta_{(a)}t)^{\frac{-2(C_{(a)}(\Lambda)+C_{(a)}(\Sigma)+C_{(a)}(\Delta))}{b_{(a)}}}. \tag{25}$$

Using the fact that quadratic Casimir invariants as well as beta functions are real, we see that magnitudes of Yukawa couplings $\hat{Y}_{\Lambda\Sigma\Delta}$s change only by $\mathcal{O}(1)$ while phases of all Yukawas do not change at all as one RG flows down from the string to the EW scale. Also, given that all four Wilson line moduli $\mathcal{A}_I$ as well as position moduli $\mathcal{Z}_I$ are stabilized at different values, we make an assumption that there will be a distinct phase factor associated with all position as well as Wilson line moduli superfields which produces an overall distinct phase factor for each possible effective Yukawa coupling corresponding to four Wilson line moduli as well as position moduli.

The other important origin of the generation of nonzero phases is given by complex soft SUSY breaking parameters $(m_i^2, \mathcal{A}_{IJK}, \mu B)$ as well supersymmetric mass term μ. The soft SUSY scalar mass terms can be made real by phase redefinition. However, in addition to the diagonal entries of sfermions corresponding to fermions with L-handed as well as R-handed chirality in the sfermion mass matrix, one gets an off-diagonal contribution because of mixing between L-R sfermion masses after EW symmetry breaking. The contribution of the same is governed by complex trilinear couplings as well as supersymmetric mass parameter μ at the EW scale. Therefore, the scalar (sfermion) fields $\tilde{f}_L$ and $\tilde{f}_R$ have been considered as linear combinations of the mass eigenstates which are obtained by diagonalizing sfermion (mass)^2 matrices [23], i.e.,

$$\tilde{f}_L = D_{f_{11}}\tilde{f}_1 + D_{f_{12}}\tilde{f}_2, \qquad \tilde{f}_R = D_{f_{21}}\tilde{f}_1 + D_{f_{22}}\tilde{f}_2, \tag{26}$$

where f corresponds to first-generation leptons and quarks and

$$D_f = \begin{pmatrix} \cos\frac{\theta_f}{2} & -\sin\frac{\theta_f}{2} e^{-i\phi_f} \\ \sin\frac{\theta_f}{2} e^{i\phi_f} & \cos\frac{\theta_f}{2} \end{pmatrix}, \tag{27}$$

and the mass matrix is given as follows:

$$M_{\tilde{f}}^2 = \begin{pmatrix} M_{\tilde{f}_L}^2 & m_u(\mathcal{A}_f^* - \mu\cot\beta) \\ m_u(\mathcal{A}_f - \mu^*\cot\beta) & M_{\tilde{f}_R}^2 \end{pmatrix}_{\rm EW}, \tag{28}$$

where $\mathcal{A}_{IJK}$ corresponds to the complex trilinear coupling parameter. Diagonalizing the above matrix by performing unitary transformation $D_f^\dagger M_{\tilde{f}}^2 d_f = {\rm diag}(M_{\tilde{f}1}^2, M_{\tilde{f}2}^2)$, where $\tan\theta_f = \frac{2|M_{\tilde{f}21}^2|}{M_{\tilde{f}11}^2 - M_{\tilde{f}22}^2}$. The eigenvalues $M_{\tilde{f}1}^2$ and $M_{\tilde{f}2}^2$ are as follows:

$$M_{\tilde{f}(1)(2)}^2 = \frac{1}{2}(M_{\tilde{f}11}^2 + M_{\tilde{f}22}^2)(+)(-) \times \frac{1}{2}[(M_{\tilde{f}11}^2 - M_{\tilde{f}22}^2)^2 + 4|M_{\tilde{f}21}^2|^2]^{\frac{1}{2}}. \tag{29}$$

For $f=e$, $\mathcal{A}_e^* = \mathcal{A}_{\mathcal{Z}_I\mathcal{A}_1\mathcal{A}_3}$; for $f=(u,d)$, $\mathcal{A}_{u/d}^* = \mathcal{A}_{\mathcal{Z}_I\mathcal{A}_2\mathcal{A}_4}$. In our model as discussed in Sec. II, we have universality in trilinear couplings with respect to position as well as Wilson line moduli. Assuming the same to be true at the EW scale, the values of the trilinear coupling parameters are $\mathcal{A}_{\mathcal{Z}_I\mathcal{A}_1\mathcal{A}_3} = \mathcal{A}_{\mathcal{Z}_I\mathcal{A}_2\mathcal{A}_4} = \mathcal{V}^{\frac{37}{36}} m_{\frac{3}{2}}$. As given in Sec. II, the value of the supersymmetric mass parameter μ at the EW scale is $\mathcal{V}^{\frac{59}{72}} m_{\frac{3}{2}}$. Also, we have universality in slepton (squark) masses of the first two generations. Therefore, $M_{\tilde{e}11}^2 = M_{\tilde{e}22}^2 = M_{\tilde{u}11}^2 = M_{\tilde{u}22}^2 \sim \mathcal{V} m_{\frac{3}{2}}^2$, and

$$|M_{\tilde{e}21}^2|^2 = m_e|\mathcal{A}_e^* - \mu\cot\beta| \equiv (\mathcal{V}^{\frac{37}{36}}) m_e m_{\frac{3}{2}} \ll M_{\tilde{e}11}^2,$$
$$|M_{\tilde{u}21}^2|^2 = m_u|\mathcal{A}_u^* - \mu\cot\beta| \equiv (\mathcal{V}^{\frac{37}{36}}) m_e m_{\frac{3}{2}} \ll M_{\tilde{u}11}^2. \tag{30}$$

Using the above, one can show that the eigenvalues of sfermion mass matrix $M_{\tilde{f}(1)(2)}^2 \sim M_{\tilde{f}_{L,R}^2} = \mathcal{V} m_{\frac{3}{2}}^2$. The aforementioned mass eigenstates can be utilized to produce a nonzero phase responsible to generate the finite EDM of the electron as well as the neutron in the one-loop diagrams involving sfermions as scalar propagators and gauginos and neutralinos as fermionic propagators.

IV. ONE-LOOP CONTRIBUTION TO THE ELECTRIC DIPOLE MOMENT

At one-loop level, for a theory of fermion ψ_f interacting with other heavy fermions' ψ_i's and heavy scalars' ϕ_k's with masses m_i, m_k and charges Q_i, Q_k, the interaction that contains CP violation, in general, is given by [23]

$$-\mathcal{L}_{\rm int} = \sum_{i,k}\bar{\psi}_f\left(K_{ik}\frac{1-\gamma_5}{2} + L_{ik}\frac{1+\gamma_5}{2}\right)\psi_i\phi_k + {\rm H.c.} \tag{31}$$

Here, $\mathcal{L}$ violates CP invariance iff ${\rm Im}(K_{ik}L^*_{ik}) \neq 0$. The one-loop EDM of the fermion f in this case is given by

$$\sum_{i,k}\frac{m_i}{(4\pi)^2 m_k^2}{\rm Im}(K_{ik}L^*_{ik})\left(Q_i A\left(\frac{m_i^2}{m_k^2}\right) + Q_k B\left(\frac{m_i^2}{m_k^2}\right)\right), \tag{32}$$

where $A(r)$ and $B(r)$ are defined by

$$A(r) = \frac{1}{2(1-r)^2}\left(3 - r + \frac{2\ln r}{1-r}\right),$$
$$B(r) = \frac{1}{2(r-1)^2}\left(1 + r + \frac{2r\ln r}{1-r}\right), \tag{33}$$

where $Q_k = Q_f - Q_i$.

We use the above-mentioned results to get an order-of-magnitude estimate of the EDM of the electron/quark in the context of $\mathcal{N} = 1$ gauged supergravity by including all SM as well as supersymmetric particles in the loop diagram. The EDMs of the neutron can be estimated by calculating the contribution of u and d quarks by using relation $d_n = (4d_d - d_u)/3$. Since in our model, we have identified both up as well as down quarks with a single Wilson line modulus, we will have the same contribution of the EDM for both up and down quarks. Hence, the neutron EDM is the same as the up-quark EDM. Therefore, in the calculations below, we will estimate the EDM of the electron and up quark only.

A. One-Loop diagrams involving neutral sfermions in the loop

1. Gaugino contribution

In this subsection, we estimate the contribution of the electron/neutron EDM at one-loop level due to the presence of a heavy gaugino nearly isospectral with heavy sfermions (for the Calabi-Yau volume $\mathcal{V} = 10^5$ in string-length units). In traditional split SUSY models discussed in the literature, the masses of sfermions are very heavy, while the masses of gauginos as well as Higginos are kept very light because of the gauge coupling unification. Therefore, one-loop diagrams involving sfermion-gaugino exchange do not give any significant contribution to the EDM of fermion. However, in the large volume $D3$-$D7$ setup that we have discussed, the gaugino as well as Higgsino also turn out to be very heavy. As it is clear from Eq. (31), the order of magnitude of the EDM at one-loop level is directly proportional to the fermion mass and inversely proportional to the sfermion masses circulating in the loop, whereas the one-loop function can almost be of $\mathcal{O}(0.1-1)$ provided either the difference between the fermion and sfermion mass is of $\mathcal{O}(1)$ or the fermion mass is very light as compared to the sfermion mass. Therefore, naively one would expect an enhancement in the order of magnitude of one-loop EDM due to the presence of heavy fermions circulating in a loop. In view of this, we estimate the contribution of the one-loop EDM of an electron as well as a neutron in the $N = 1$ gauged supergravity limit of large volume $D3/D7$ μ-split SUSY model discussed in Sec. II. However, the CP violation (imaginary phases) can be induced in a loop diagram by considering diagonalized eigenstates of sfermion mass matrix as propagators in the loop. The loop diagram is given in Fig. 1. The effective one-loop operator given in Eq. (31) can be recast in the following form:

$$\mathcal{L}_{\rm int} = \sum_{i=e,u,d}\bar{\psi}_{f_i}\left(K_i\frac{1-\gamma_5}{2} + L_i\frac{1+\gamma_5}{2}\right)\phi_{\tilde{f}_i}\tilde{\lambda}^0 + {\rm H.c.} \tag{34}$$

For $i = 1, 2$, the above equation can be expanded as

$$-\mathcal{L}_{\rm int} = \bar{\psi}_f\left(K_1\frac{1-\gamma_5}{2} + L_1\frac{1+\gamma_5}{2}\right)\phi_{\tilde{f}_1}\tilde{\lambda}^0 + \bar{\psi}_f\left(K_2\frac{1-\gamma_5}{2} + L_2\frac{1+\gamma_5}{2}\right)\phi_{\tilde{f}_2}\tilde{\lambda}^0 + {\rm H.c.}, \tag{35}$$

and the one-loop EDM of the fermion f in this case will be given as

$$\frac{m_{\tilde{\lambda}^0_i}}{(4\pi)^2}\left[\frac{1}{m_{\tilde{f}_1}{}^2}{\rm Im}(K_1L_1^*)\left(Q'_{\tilde{f}_1}B\left(\frac{m^2_{\tilde{\lambda}^0}}{m_{\tilde{f}_1}{}^2}\right)\right) + \frac{1}{m_{\tilde{f}_2}{}^2}{\rm Im}(K_2L_2^*)\left(Q_{\tilde{f}_2}B\left(\frac{m^2_{\tilde{\lambda}^0}}{m_{\tilde{f}_2}{}^2}\right)\right)\right], \tag{36}$$

where $m_{\tilde{\lambda}^0}$ corresponds to the gaugino mass, $m_{\tilde{f}_1}$ and $m_{\tilde{f}_2}$ correspond to the masses of the eigenstates of the diagonalized sfermion mass matrix, and $Q'_{\tilde{f}_i}$ corresponds to effective charge defined as $Q'_{\tilde{f}_i} \sim Q_{\tilde{f}_i}C_{\tilde{f}_i\tilde{f}_i\gamma}$, where $C_{\tilde{f}_i\tilde{f}_i\gamma}$

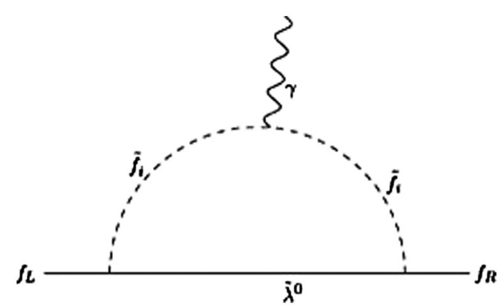

FIG. 1. One-loop diagram involving gauginos.

will be the volume-suppression factor coming from the sfermion-photon-sfermion vertex.

To determine the value of the one-loop EDM in this case, we first calculate the contribution of the required vertices involved in Fig. 1. In $\mathcal{N}=1$ gauged supergravity, lepton (quark)-slepton(squark)-gaugino interaction ia governed by [39] the following term:

$$\mathcal{L}_{f-\tilde{f}-\tilde{\lambda}^0} = g_{YM} g_{J\bar{T}_B} X^{*B} \bar{\chi}^{\bar{J}} \tilde{\lambda}^0 + \partial_J T_B D^B \bar{\chi}^{\bar{J}} \lambda^0,$$

where $\bar{\chi}^{\bar{J}}$ corresponds to a spin-$\frac{1}{2}$ fermion, X^{*B} is the Killing isometry vector, and λ^0 corresponds to an $SU(2)$-singlet component of the neutral gaugino. Though the gauge coupling g_{YM} is real, the nonzero phase factor is produced from the moduli space metric component $g_{J\bar{T}_B}$ and is associated with the volume-suppression factor arising from the same. Hence, the effective gauge coupling interaction vertex generates a particular phase factor which we consider to be of the order $\mathcal{O}(1)$.

We repeatedly mention that to get the numerical estimate of the contribution of the aforementioned vertices, we use the identification described in [38] according to which fermionic superpartners of $\mathcal{A}_1$ and $\mathcal{A}_3$ can be identified, respectively, with e_L and e_R, and the fermionic superpartners of $\mathcal{A}_2$ and $\mathcal{A}_4$ can be identified, respectively, with the first-generation quarks: u_L/d_L and u_R/d_R. In principle, an incoming left-handed electron (quark) can couple with scalar superpartners of both left-handed as well as right-handed leptons (quarks). Therefore, for a left-handed electron e_L interacting with a slepton as well as gaugino, the interaction vertex will be given as

$$\mathcal{L}_{e_L-\tilde{e}-\tilde{\Lambda}^0} = g_{YM} g_{\mathcal{A}_1\bar{T}_B} X^{*B} \bar{\chi}^{\bar{\mathcal{A}}_1} \tilde{\lambda}^0 + \partial_{a_1} T_B D^B \bar{\chi}^{\bar{a}_1} \lambda^0. \quad (37)$$

To calculate the contribution of the $e_L - \tilde{e_L} - \tilde{\lambda}^0$ vertex, we expand $g_{\mathcal{A}_1\bar{T}_B}$ in the fluctuations linear in $\mathcal{A}_1$ around its stabilized VEV. In terms of an undiagonalized basis, we have $g_{T_B\bar{a}_1} \to -\mathcal{V}^{-\frac{1}{4}}(a_1 - \mathcal{V}^{-\frac{2}{9}})$. Using $T_B = \mathrm{vol}(\sigma_B) - C_{I\bar{J}} a_I \bar{a}_{\bar{J}} + \mathrm{H.c.}$, where the values of intersection matrices $C_{I\bar{J}}$ are given in the Appendix of [38]. Utilizing those values, we get $\partial_{a_1} T_B \to \mathcal{V}^{\frac{10}{9}}(a_1 - \mathcal{V}^{-\frac{2}{9}})$. Using the argument that $g_{YM} g_{T_B\bar{a}_1} \sim \mathcal{O}(1) g_{YM} g_{T_B\bar{\mathcal{A}}_1}$ as shown in [38], incorporating the values of $X^B = -6i\kappa_4^2\mu_7 Q_{T_B}$, $\kappa_4^2\mu_7 \sim \frac{1}{\mathcal{V}}$, $D^B = \frac{4\pi\alpha'\kappa_4^2\mu_7 Q_B v^B}{\mathcal{V}}$, and $Q_{T_B} \sim \mathcal{V}^{\frac{1}{3}}(2\pi\alpha')^2 \tilde{f}$, we get the contribution of the physical gaugino($\tilde{\lambda}^0$)-lepton(e_L)-slepton($\tilde{e_L}$) interaction vertex as follows:

$$|C_{e_L\tilde{e_L}\tilde{\lambda}^0}| \equiv \frac{\mathcal{V}^{-\frac{2}{9}}\tilde{f}}{\sqrt{\hat{K}_{\mathcal{A}_1\bar{\mathcal{A}}_1}}\sqrt{\hat{K}_{\mathcal{A}_1\bar{\mathcal{A}}_1}}} \tilde{\mathcal{A}}_1 \bar{\chi}^{\bar{\mathcal{A}}_1} \tilde{\lambda}^0 \equiv \tilde{f}(\mathcal{V}^{-1}), \quad (38)$$

where $\tilde{f}$ is the dilute flux, and the upper limit of the same as calculated in [37] is $\mathcal{V}^{-\frac{23}{30}} \sim \mathcal{O}(10^{-4})$ for Calabi-Yau volume $\mathcal{V} \sim 10^5$.

Similarly, the contribution of the physical gaugino($\tilde{\lambda}^0$)-quark(u_L)-squark($\tilde{u_L}$) interaction vertex will be given by expanding $g_{\mathcal{A}_2\bar{T}_B}$ in the fluctuations linear in $\mathcal{A}_2$ around its stabilized VEV. Doing so, one will get $g_{T_B\bar{a}_2} \to -\mathcal{V}^{-\frac{5}{4}}(a_2 - \mathcal{V}^{-\frac{1}{3}})$, $\partial_{a_2} T_B \to \mathcal{V}^{\frac{1}{9}}(a_2 - \mathcal{V}^{-\frac{1}{3}})$, and

$$|C_{u_L\tilde{u_L}\tilde{\lambda}^0}| \equiv \frac{\mathcal{V}^{-\frac{11}{9}}\tilde{f}}{\sqrt{\hat{K}_{\mathcal{A}_2\bar{\mathcal{A}}_2}}\sqrt{\hat{K}_{\mathcal{A}_2\bar{\mathcal{A}}_2}}} \tilde{\mathcal{A}}_2 \bar{\chi}^{\bar{\mathcal{A}}_2} \tilde{\lambda}^0 \equiv \tilde{f}(\mathcal{V}^{-\frac{4}{5}}). \quad (39)$$

The gaugino($\tilde{\lambda}^0$)-fermion(f_L)-sfermion($\tilde{f}_R$) vertex does not possess $SU(2)$ electroweak symmetry. However, the terms in the supergravity Lagrangian preserve $SU(2)$ EW symmetry. Therefore, we first generate a term of the type $f_L \tilde{f}_R \tilde{\lambda}^0 H_L$ wherein H_L is an $SU(2)_L$ Higgs doublet. After spontaneous breaking of the EW symmetry when H_L acquires a nonzero VEV $\langle H^0 \rangle$, this term generates $\langle H^0 \rangle f_L \tilde{f}_R \tilde{\lambda}^0$. For $f_{L,R} = e_{L,R}$, by expanding $g_{\mathcal{A}_1\bar{T}_B}$ in the fluctuations linear in $\mathcal{Z}_i$ and then linear in $\mathcal{A}_3$ around their stabilized value, we have $g_{T_B\bar{\mathcal{A}}_1} \to \mathcal{V}^{-\frac{13}{36}}\langle \mathcal{Z}_1 \rangle(\mathcal{A}_3 - \mathcal{V}^{-\frac{13}{18}})$. The contribution of physical gaugino($\tilde{\lambda}^0$)-lepton(e_L)-slepton($\tilde{e_R}$) interaction vertex will be as follows:

$$|C_{e_L\tilde{e_R}\tilde{\lambda}^0}| \equiv \frac{g_{YM} g_{T_B\bar{\mathcal{A}}_1} X^{T_B} \sim \mathcal{V}^{-\frac{19}{18}}\tilde{f}}{\sqrt{\hat{K}_{\mathcal{Z}_1\bar{\mathcal{Z}}_1}\hat{K}_{\mathcal{A}_1\bar{\mathcal{A}}_1}\hat{K}_{\mathcal{A}_3\bar{\mathcal{A}}_3}}} \tilde{\mathcal{A}}_3 \bar{\chi}^{\bar{\mathcal{A}}_1} \tilde{\lambda}^0 \equiv \tilde{f}(\mathcal{V}^{-\frac{15}{9}}). \quad (40)$$

For $f_{L,R} = u_{L,R}$, by expanding $g_{\mathcal{A}_2\bar{T}_B}$ in the fluctuations linear in $\mathcal{Z}_i$ and then linear in $\mathcal{A}_4$ around their stabilized value, we have $g_{T_B\bar{\mathcal{A}}_2} \to \mathcal{V}^{-\frac{13}{36}}\langle \mathcal{Z}_1 \rangle(\mathcal{A}_4 - \mathcal{V}^{-\frac{11}{9}})$ and

$$|C_{u_L\tilde{u_R}\tilde{\lambda}^0}| \equiv \frac{g_{YM} g_{T_B\bar{\mathcal{A}}_2} X^{T_B} \sim \mathcal{V}^{-\frac{19}{18}}\tilde{f}}{\sqrt{\hat{K}_{\mathcal{Z}_1\bar{\mathcal{Z}}_1}\hat{K}_{\mathcal{A}_2\bar{\mathcal{A}}_2}\hat{K}_{\mathcal{A}_4\bar{\mathcal{A}}_4}}} \tilde{\mathcal{A}}_4 \bar{\chi}^{\bar{\mathcal{A}}_2} \tilde{\lambda}^0 \equiv \tilde{f}(\mathcal{V}^{-\frac{14}{9}}). \quad (41)$$

Similarly, the outgoing right-handed electron (quark) can couple with both the left-handed as well as right-handed sleptons (squarks) and include the gaugino($\tilde{\lambda}^0$)-fermion(f_R)-sfermion($\tilde{f}_L$) vertex in a loop diagram. The same does not possess $SU(2)$ EW symmetry. For $f_{L,R} = e_{L,R}$, by expanding $g_{\mathcal{A}_3\bar{T}_B}$ first in the fluctuations linear in $\mathcal{Z}_1$ and then linear in $\mathcal{A}_1$ around their stabilized VEVs, we have $g_{T_B\bar{\mathcal{A}}_3} \to -\mathcal{V}^{-\frac{13}{36}}\langle \mathcal{Z}_1 \rangle(\mathcal{A}_1 - \mathcal{V}^{-\frac{2}{9}})$. The physical gaugino($\tilde{\lambda}^0$)-lepton($e_R$)-slepton($\tilde{e_L}$) interaction vertex will be given as

$$|C_{e_R\tilde{e_L}\tilde{\lambda}^0}| \equiv \frac{g_{YM}g_{T_B\bar{\mathcal{A}}_3}X^{T_B}\sim\mathcal{V}^{-\frac{19}{18}}\tilde{f}}{\sqrt{\hat{K}_{\mathcal{Z}_1\bar{\mathcal{Z}}_1}\hat{K}_{\mathcal{A}_1\bar{\mathcal{A}}_1}\hat{K}_{\mathcal{A}_3\bar{\mathcal{A}}_3}}}\tilde{\mathcal{A}}_1\bar{\chi}^{\bar{\mathcal{A}}_3}\tilde{\lambda}^0 \equiv \tilde{f}(\mathcal{V}^{-\frac{15}{9}}). \tag{42}$$

For $f_{L,R}=u_{L,R}$, one gets $g_{T_B\bar{\mathcal{A}}_4}\to -\mathcal{V}^{-\frac{13}{36}}\langle\mathcal{Z}_1\rangle(\mathcal{A}_2-\mathcal{V}^{-\frac{1}{3}})$ and

$$|C_{u_R\tilde{u_L}\tilde{\lambda}^0}| \equiv \frac{g_{YM}g_{T_B\bar{\mathcal{A}}_2}X^{T_B}\sim\mathcal{V}^{-\frac{19}{18}}\tilde{f}}{\sqrt{\hat{K}_{\mathcal{Z}_1\bar{\mathcal{Z}}_1}\hat{K}_{\mathcal{A}_2\bar{\mathcal{A}}_2}\hat{K}_{\mathcal{A}_4\bar{\mathcal{A}}_4}}}\tilde{\mathcal{A}}_2\bar{\chi}^{\bar{\mathcal{A}}_4}\tilde{\lambda}^0 \equiv \tilde{f}(\mathcal{V}^{-\frac{14}{9}}). \tag{43}$$

To calculate the contribution of the $e_R-\tilde{e_R}-\tilde{\lambda}^0$ vertex, we expand $g_{\mathcal{A}_3\bar{T}_B}$ in the fluctuations linear in $\mathcal{A}_3$ and obtain $g_{T_B\bar{\mathcal{A}}_3}\to -\mathcal{V}^{\frac{7}{9}}(\mathcal{A}_3-\mathcal{V}^{\frac{13}{18}})$, $\partial_{\mathcal{A}_3}T_B\to\mathcal{V}^{\frac{19}{9}}(\mathcal{A}_3-\mathcal{V}^{-\frac{13}{18}})$. Utilizing this, the physical gaugino($\tilde{\lambda}^0$)-lepton(e_R)-slepton($\tilde{e_R}$) interaction vertex will be given as

$$|C_{e_R\tilde{e_R}\tilde{\lambda}^0}| \equiv \frac{\mathcal{V}^{\frac{7}{9}}\tilde{f}}{\sqrt{\hat{K}_{\mathcal{A}_3\bar{\mathcal{A}}_3}\hat{K}_{\mathcal{A}_3\bar{\mathcal{A}}_3}}}\tilde{\mathcal{A}}_3\bar{\chi}^{\bar{\mathcal{A}}_3}\tilde{\lambda}^0 \equiv \tilde{f}(\mathcal{V}^{-\frac{3}{5}}). \tag{44}$$

Similarly, by expanding $g_{\mathcal{A}_4\bar{T}_B}$ in the fluctuations linear in $\mathcal{A}_4$, we will have $g_{T_B\bar{\mathcal{A}}_4}\to -\mathcal{V}^{\frac{16}{9}}(\mathcal{A}_4-\mathcal{V}^{\frac{11}{9}})$, $\partial_{\mathcal{A}_4}T_B\to\mathcal{V}^{\frac{28}{9}}(\mathcal{A}_4-\mathcal{V}^{-\frac{11}{9}})$, and

$$|C_{u_R\tilde{u_R}\tilde{\lambda}^0}| \equiv \frac{\mathcal{V}^{\frac{16}{9}}\tilde{f}}{\sqrt{\hat{K}_{\mathcal{A}_4\bar{\mathcal{A}}_4}\hat{K}_{\mathcal{A}_4\bar{\mathcal{A}}_4}}}\tilde{\mathcal{A}}_4\bar{\chi}^{\bar{\mathcal{A}}_4}\tilde{\lambda}^0 \equiv \tilde{f}(\mathcal{V}^{-\frac{3}{5}}). \tag{45}$$

To determine the contribution of effective charge Q_i', we need to evaluate the contribution of sfermion($\tilde{f}_i$)-photon(γ)-sfermion($\tilde{f}_i$) vertices which are expressed in terms of $\tilde{f}_{L/R}$ basis as below

$$\begin{aligned}C_{\tilde{f}_1\tilde{f}_1\gamma} &\sim D_{f_{11}}D^*_{f_{11}}C_{\tilde{f}_L\tilde{f}_L{}^*\gamma}+(D_{f_{11}}D^*_{f_{12}}+D_{f_{12}}D^*_{f_{12}})C_{\tilde{f}_L\tilde{f}_R{}^*\gamma}\\&\quad+D_{f_{12}}D^*_{f_{12}}C_{\tilde{f}_R\tilde{f}_R{}^*\gamma},\\ C_{\tilde{f}_2\tilde{f}_2\gamma} &\sim D_{f_{21}}D^*_{f_{21}}C_{\tilde{f}_L\tilde{f}_L{}^*\gamma}+(D_{f_{21}}D^*_{f_{22}}+D_{f_{22}}D^*_{f_{21}})C_{\tilde{f}_L\tilde{f}_R{}^*\gamma}\\&\quad+D_{f_{22}}D^*_{f_{22}}C_{\tilde{f}_R\tilde{f}_R{}^*\gamma}.\end{aligned} \tag{46}$$

The sfermion-sfermion-photon vertex can be evaluated from the bulk kinetic term in the $\mathcal{N}=1$ gauged supergravity action as given below

$$\mathcal{L}=\frac{1}{\kappa_4^2\mathcal{V}^2}G^{T_B\bar{T}_B}\tilde{\nabla}_\mu T_B\tilde{\nabla}^\mu\bar{T}_{\bar{B}}, \tag{47}$$

where

$$\begin{aligned}&\tilde{\nabla}_\mu T_B=\partial_\mu T_B+6i\kappa_4^2\mu_7 lQ_{T_B}A_\mu;\\ &T_B\sim\sigma_B+\left(i\kappa_{Bbc}c^b\mathcal{B}^c+\kappa_B+\frac{i}{(\tau-\bar{\tau})}\kappa_{Bbc}\mathcal{G}^b(\mathcal{G}^c-\bar{\mathcal{G}}^c)i\delta^B_B\kappa_4^2\mu_7 l^2C^{I\bar{J}}_B a_I\bar{a}_{\bar{J}}+\frac{3i}{4}\delta^B_B\tau Q_{\tilde{f}}+i\mu_3 l^2(\omega_B)_{i\bar{j}}z^i\left(\bar{z}^{\bar{j}}-\frac{i}{2}\bar{z}^{\tilde{a}}(\bar{\mathcal{P}}_{\tilde{a}})^{\bar{j}}_l z^l\right)\right).\end{aligned} \tag{48}$$

The form of expression that eventually leads to give the contribution of required sfermion-sfermion-photon vertex is given below

$$C_{f_{L/R}f^*_{L/R}\gamma}\sim\frac{6i\kappa_4^2\mu_7 2\pi\alpha' Q_B G^{T_B\bar{T}_B}}{\kappa_4^2\mathcal{V}^2}\times A^\mu\partial_\mu(\kappa_4^2\mu_7(2\pi\alpha')^2C_{i\bar{j}}\mathcal{A}_i\bar{\mathcal{A}}_{\bar{j}}). \tag{49}$$

Using $G^{T_B\bar{T}_B}(\mathrm{EW})\sim\mathcal{V}^{\frac{7}{3}}$ [the large value is justified by obtaining the $\mathcal{O}(1)$ SM fermion-fermion-photon coupling vertex in $\mathcal{N}=1$ gauged supergravity action; see details therein], $Q_B\sim\mathcal{V}^{\frac{1}{3}}\tilde{f}$, $\kappa_4^2\mu_7\sim\frac{1}{\mathcal{V}}$, the above expression reduces to $|C_{f_{L/R}f^*_{L/R}\gamma}|\equiv\mathcal{V}^{\frac{1}{3}}A^\mu\partial_\mu(\kappa_4^2\mu_7(2\pi\alpha')^2C_{i\bar{j}}\tilde{\mathcal{A}}_i\tilde{\mathcal{A}}_{\bar{j}})$. For $i=j=1$, $\kappa_4^2\mu_7(2\pi\alpha')^2C_{1\bar{1}}\sim\mathcal{V}^{\frac{10}{9}}$ as given in the Appendix of [38]. Using the same,

$$|C_{\tilde{e_L}\tilde{e_L}^*\gamma}|\equiv\frac{\mathcal{V}^{\frac{16}{9}}\tilde{f}}{\sqrt{\hat{K}_{\mathcal{A}_1\bar{\mathcal{A}}_1}\hat{K}_{\mathcal{A}_1\bar{\mathcal{A}}_1}}}\equiv(\mathcal{V}^{\frac{44}{45}}\tilde{f})\tilde{\mathcal{A}}_1A^\mu\partial_\mu\tilde{\mathcal{A}}_1. \tag{50}$$

For $i=1$, $j=3$; $\kappa_4^2\mu_7(2\pi\alpha')^2C_{1\bar{3}}\sim\mathcal{V}^{\frac{29}{18}}$, we have

$$|C_{\tilde{e_L}\tilde{e_R}^*\gamma}|\equiv\frac{\mathcal{V}^{\frac{41}{18}}\tilde{f}}{\sqrt{\hat{K}_{\mathcal{A}_1\bar{\mathcal{A}}_1}\hat{K}_{\mathcal{A}_3\bar{\mathcal{A}}_3}}}\equiv(\mathcal{V}^{\frac{53}{45}}\tilde{f})\tilde{\mathcal{A}}_1A^\mu\partial_\mu\tilde{\mathcal{A}}_3. \tag{51}$$

For $i=j=3$; $\kappa_4^2\mu_7(2\pi\alpha')^2C_{3\bar{3}}\sim\mathcal{V}^{\frac{19}{9}}$ and

$$|C_{\tilde{e_R}\tilde{e_R}^*\gamma}|\equiv\frac{\mathcal{V}^{\frac{25}{9}}\tilde{f}}{\sqrt{\hat{K}_{\mathcal{A}_3\bar{\mathcal{A}}_3}\hat{K}_{\mathcal{A}_3\bar{\mathcal{A}}_3}}}\equiv(\tilde{f}\mathcal{V}^{\frac{62}{45}})\tilde{\mathcal{A}}_3A^\mu\partial_\mu\tilde{\mathcal{A}}_3. \tag{52}$$

For $i=j=2$, $\kappa_4^2\mu_7(2\pi\alpha')^2C_{2\bar{2}}\sim\mathcal{V}^{\frac{1}{9}}$ and

$$|C_{\tilde{u}_L\tilde{u}_L{}^*\gamma}| \equiv \frac{\mathcal{V}^{\frac{7}{9}}\tilde{f}}{\sqrt{\hat{K}_{\mathcal{A}_2\bar{\mathcal{A}}_2}\hat{K}_{\mathcal{A}_2\bar{\mathcal{A}}_2}}} \equiv (\tilde{f}\mathcal{V}^{\frac{53}{45}})\tilde{\mathcal{A}}_2 A^\mu \partial_\mu \tilde{\mathcal{A}}_2. \quad (53)$$

For $i = 2$, $j = 4$, $\kappa_4^2\mu_7(2\pi\alpha')^2 C_{2\bar{4}} \sim \mathcal{V}^{\frac{29}{18}}$ and

$$|C_{\tilde{u}_L\tilde{u}_R{}^*\gamma}| \equiv \frac{\mathcal{V}^{\frac{41}{18}}\tilde{f}}{\sqrt{\hat{K}_{\mathcal{A}_2\bar{\mathcal{A}}_2}\hat{K}_{\mathcal{A}_4\bar{\mathcal{A}}_4}}} \equiv (\tilde{f}\mathcal{V}^{\frac{23}{18}})\tilde{\mathcal{A}}_2 A^\mu \partial_\mu \tilde{\mathcal{A}}_4. \quad (54)$$

For $i = j = 4$, $\kappa_4^2\mu_7(2\pi\alpha')^2 C_{4\bar{4}} \sim \mathcal{V}^{\frac{28}{9}}$ and

$$|C_{\tilde{u}_R\tilde{u}_R{}^*\gamma}| \equiv \frac{\mathcal{V}^{\frac{34}{9}}\tilde{f}}{\sqrt{\hat{K}_{\mathcal{A}_4\bar{\mathcal{A}}_4}\hat{K}_{\mathcal{A}_4\bar{\mathcal{A}}_4}}} \equiv (\tilde{f}\mathcal{V}^{\frac{62}{45}})\tilde{\mathcal{A}}_4 A^\mu \partial_\mu \tilde{\mathcal{A}}_4. \quad (55)$$

Substituting the results given in Eqs. (50)–(55) in Eq. (46), the volume-suppression factors corresponding to scalar-scalar-photon vertices are given as follows:

$$\begin{aligned}
C_{\tilde{e}_1\tilde{e}_1\gamma} &\equiv \tilde{f}(\mathcal{V}^{\frac{44}{45}}\cos^2\theta_e - \mathcal{V}^{\frac{53}{45}}\cos\theta_e\sin\theta_e(e^{i\phi_e} + e^{-i\phi_e})e^{i\phi_{g_e}} + \mathcal{V}^{\frac{62}{45}}\sin^2\theta_e),\\
C_{\tilde{e}_2\tilde{e}_2\gamma} &\equiv \tilde{f}(\mathcal{V}^{\frac{44}{45}}\sin^2\theta_e + \mathcal{V}^{\frac{53}{45}}\cos\theta_e\sin\theta_e(e^{i\phi_e} + e^{-i\phi_e}))e^{-i\phi_{g_e}} + \mathcal{V}^{\frac{62}{45}}\cos^2\theta_e),\\
C_{\tilde{u}_1\tilde{u}_1\gamma} &\equiv \tilde{f}(\mathcal{V}^{\frac{53}{45}}\cos^2\theta_u - \mathcal{V}^{\frac{23}{18}}\cos\theta_u\sin\theta_e(e^{i\phi_u} + e^{-i\phi_u}))e^{i\phi_{g_u}} + \mathcal{V}^{\frac{62}{45}}\sin^2\theta_u),\\
C_{\tilde{u}_2\tilde{u}_2\gamma} &\equiv \tilde{f}(\mathcal{V}^{\frac{53}{45}}\sin^2\theta_u + \mathcal{V}^{\frac{23}{18}}\cos\theta_u\sin\theta_u(e^{i\phi_u} + e^{-i\phi_u}))e^{-i\phi_{g_u}} + \mathcal{V}^{\frac{62}{45}}\cos^2\theta_u),
\end{aligned} \quad (56)$$

where ϕ_{g_e} and ϕ_{g_u} are phase factors associated with $C_{\tilde{e}_L\tilde{e}_R^*\gamma}$ and $C_{\tilde{u}_L\tilde{u}_R^*\gamma}$ [we consider the same to be $\mathcal{O}(1)$]. Now, the Lagrangian relevant to the couplings involved in the one-loop diagram shown in Fig. 1 is given as

$$\mathcal{L} = C_{f_L\tilde{f}_L^*\tilde{\lambda}_i^0} f_L\tilde{f}_L\tilde{\lambda}^0 + C_{f_L\tilde{f}_R^*\tilde{\lambda}_i^0} f_L\tilde{f}_R\tilde{\lambda}^0 + C_{f_R^*\tilde{f}_L\tilde{\lambda}_i^0} f_R\tilde{f}_L\tilde{\lambda}^0 + C_{f_R^*\tilde{f}_R\tilde{\lambda}_i^0} f_R\tilde{f}_R\tilde{\lambda}^0, \quad (57)$$

where, from Eqs. (38)–(45), we have

$$\begin{aligned}
&|C_{e_L\tilde{e}_L^*\tilde{\lambda}_i^0}| \equiv \tilde{f}\mathcal{V}^{-1}, \qquad |C_{e_R\tilde{e}_R^*\tilde{\lambda}_i^0}| \equiv \tilde{f}\mathcal{V}^{-\frac{3}{5}}, \qquad |C_{e_R^*\tilde{e}_L\tilde{\lambda}_i^0}| \equiv |C_{e_L^*\tilde{e}_R\tilde{\lambda}_i^0}| \equiv \tilde{f}\mathcal{V}^{-\frac{15}{9}},\\
&|C_{u_L\tilde{u}_L^*\tilde{\lambda}_i^0}| \equiv \tilde{f}\mathcal{V}^{-\frac{4}{5}}, \qquad |C_{u_R\tilde{u}_R^*\tilde{\lambda}_i^0}| \equiv \tilde{f}\mathcal{V}^{-\frac{3}{5}}, \qquad |C_{u_R^*\tilde{u}_L\tilde{\lambda}_i^0}| \equiv |C_{u_L^*\tilde{u}_R\tilde{\lambda}_i^0}| \equiv \tilde{f}\mathcal{V}^{-\frac{14}{9}}.
\end{aligned} \quad (58)$$

Writing $\tilde{f}_L$ as well as $\tilde{f}_R$ given in Eq. (57) in terms of diagonalized basis $\tilde{f}_1$ and $\tilde{f}_2$, the equation takes the form of Eq. (35):

$$\begin{aligned}
\mathcal{L}_{\rm int} &= \bar{\chi}_f\left((C_{\lambda_i^0 f_L\tilde{f}_L}D_{f_{11}} + C_{\lambda_i^0 f_L\tilde{f}_R}D_{f_{21}})\frac{1+\gamma_5}{2} + (C_{\lambda_i^0 f_R\tilde{f}_L}D_{f_{11}} + C_{\lambda_i^0 f_R\tilde{f}_R}D_{f_{21}})\frac{1-\gamma_5}{2}\right)\phi_{f_1}\tilde{\lambda}^0\\
&\quad + \bar{\chi}_f\left((C_{\lambda_i^0 f_L\tilde{f}_L}D_{f_{12}} + C^{\chi_i^0 f_L\tilde{f}_R}D_{f_{22}})\frac{1+\gamma_5}{2} + (C_{\lambda_i^0 f_R\tilde{f}_L}D_{f_{12}} + C_{\lambda_i^0 f_R\tilde{f}_R}D_{f_{22}})\frac{1-\gamma_5}{2}\right)\phi_{f_2}\tilde{\lambda}^0 + \text{H.c.}
\end{aligned} \quad (59)$$

Using Eq. (36), the EDM expression will take the form

$$\begin{aligned}
\left.\frac{d_f}{e}\right|_{\lambda_i^0} &= \frac{m_{\tilde{\lambda}_i^0}}{(4\pi)^2}\left[\frac{1}{m^2_{\tilde{f}_1}}\text{Im}(C_{\lambda_i^0 f_L\tilde{f}_L}C_{\lambda_i^0 f_R\tilde{f}_R}D_{f_{11}}D^*_{f_{21}} + C_{\lambda_i^0 f_L\tilde{f}_R}C_{\chi_i^0 f_R\tilde{f}_L}D_{f_{21}}D^*_{f_{11}})Q'_{\tilde{f}_1}B\left(\frac{m^2_{\tilde{\lambda}_i^0}}{m_{\tilde{f}_1}{}^2}\right)\right.\\
&\quad \left.+\frac{1}{m^2_{\tilde{f}_2}}\text{Im}(C_{\lambda_i^0 f_L\tilde{f}_L}C_{\chi_i^0 f_R\tilde{f}_R}D_{f_{12}}D^*_{f_{22}} + C_{\lambda_i^0 f_L\tilde{f}_R}C_{\lambda_i^0 f_R\tilde{f}_L}D_{f_{22}}D^*_{f_{12}})Q'_{\tilde{f}_2}B\left(\frac{m^2_{\tilde{\lambda}_i^0}}{m_{\tilde{f}_2}{}^2}\right)\right].
\end{aligned} \quad (60)$$

Considering $f_{L,R} = e_{L,R}$, incorporating the results of the interaction vertices as given in Eq. (58), and using the assumption that the phase factors associated with effective gauge couplings are $\mathcal{O}(1)$, the dominant contribution of the electron EDM is given as

$$\left.\frac{d_e}{e}\right|_{\lambda_i^0} \equiv \frac{m_{\tilde{\lambda}^0}(\tilde{f}^2\mathcal{V}^{-\frac{8}{5}}\sin\theta_e\sin\phi_e)}{(4\pi)^2}\left[\frac{C_{\tilde{e}_2\tilde{e}_2{}^*\gamma}}{m^2_{\tilde{e}_2}}B\left(\frac{m^2_{\tilde{\lambda}^0}}{m_{\tilde{e}_2}{}^2}\right) - \frac{C_{\tilde{e}_1\tilde{e}_1{}^*\gamma}}{m^2_{\tilde{e}_1}}B\left(\frac{m^2_{\tilde{\lambda}^0}}{m_{\tilde{e}_1}{}^2}\right)\right]. \quad (61)$$

For $f_{L,R} = u_{L,R}$, the quark EDM will be given as

$$\frac{d_u}{e}\bigg|_{\lambda_i^0} \equiv \frac{m_{\tilde{\lambda}^0}(\tilde{f}^2\mathcal{V}^{-\frac{7}{5}}\sin\theta_u\sin\phi_u)}{(4\pi)^2}\left[\frac{C_{\tilde{u}_2\tilde{u}_2^*\gamma}}{m^2_{\tilde{u}_2}}B\left(\frac{m^2_{\tilde{\lambda}^0}}{m_{\tilde{u}_2}{}^2}\right) - \frac{C_{\tilde{u}_1\tilde{u}_1^*\gamma}}{m^2_{\tilde{u}_1}}B\left(\frac{m^2_{\tilde{\lambda}^0}}{m_{\tilde{u}_1}{}^2}\right)\right]. \tag{62}$$

Putting the values[3] of $C_{\tilde{e}_i\tilde{e}_i^*\gamma}$ and $C_{\tilde{u}_i\tilde{u}_i^*\gamma}$ as given in Eq. (56), we get

$$\frac{d_e}{e}\bigg|_{\lambda^0} \equiv \frac{m_{\tilde{\lambda}^0}(\tilde{f}^2\mathcal{V}^{-\frac{8}{5}}\sin\theta_e\sin\phi_e)}{(4\pi)^2}\mathcal{V}^{\frac{62}{45}}\tilde{f}\left[\frac{\cos^2\theta_e}{m^2_{\tilde{e}_2}}B\left(\frac{m^2_{\tilde{\lambda}^0}}{m_{\tilde{e}_2}{}^2}\right) - \frac{\sin^2\theta_e}{m^2_{\tilde{e}_1}}B\left(\frac{m^2_{\tilde{\lambda}^0}}{m_{\tilde{e}_1}{}^2}\right)\right],$$
$$\frac{d_u}{e}\bigg|_{\lambda^0} \equiv \frac{m_{\tilde{\lambda}^0}(\tilde{f}^2\mathcal{V}^{-\frac{7}{5}}\sin\theta_u\sin\phi_u)}{(4\pi)^2}\mathcal{V}^{\frac{62}{45}}\tilde{f}\left[\frac{\cos^2\theta_u}{m^2_{\tilde{u}_2}}B\left(\frac{m^2_{\tilde{\lambda}^0}}{m_{\tilde{u}_2}{}^2}\right) - \frac{\sin^2\theta_u}{m^2_{\tilde{u}_1}}B\left(\frac{m^2_{\tilde{\lambda}^0}}{m_{\tilde{u}_1}{}^2}\right)\right]. \tag{63}$$

Here, $\sin\theta_f = \frac{2|M_{\tilde{f}^2_{21}}|}{\sqrt{(M^2_{\tilde{f}_L}-M^2_{\tilde{f}_R})^2+4M^4_{\tilde{f}21}}}$. As discussed in Sec. II, in our model, we have $M^2_{\tilde{e}_L}=M^2_{\tilde{e}_R}=M^2_{\tilde{u}_L}=M^2_{\tilde{u}_R}\sim\mathcal{V}m^2_{\frac{3}{2}}$. Using the same, we get $\sin\theta_e=\sin\theta_u=1$. Also, we assume $\phi_{e,u}=(0,\frac{\pi}{2}]$. As explained in Sec. III, $m^2_{\tilde{f}_1}=m^2_{\tilde{f}_2}=m^2_{\tilde{f}_L/\tilde{f}_R}=\mathcal{V}m^2_{\frac{3}{2}}$. Utilizing the same and the value of the gaugino mass $m^2_{\tilde{\lambda}^0}=\mathcal{V}^{\frac{4}{3}}m^2_{\frac{3}{2}}$, we get

$$B\left(\frac{m^2_{\tilde{\lambda}^0}}{m_{\tilde{f}_i}{}^2}\right) = \frac{1}{2\left(\frac{m^2_{\tilde{\lambda}^0}}{m_{\tilde{f}_i}{}^2}-1\right)^2}\left(1+\frac{m^2_{\tilde{\lambda}^0}}{m_{\tilde{f}_i}{}^2}+\frac{2\frac{m^2_{\tilde{\lambda}^0}}{m_{\tilde{f}_i}{}^2}\ln\left(\frac{m^2_{\tilde{\lambda}^0}}{m_{\tilde{f}_i}{}^2}\right)}{1-\frac{m^2_{\tilde{\lambda}^0}}{m_{\tilde{f}_i}{}^2}}\right)$$
$$\sim\frac{m_{\tilde{f}_i}{}^2}{m^2_{\tilde{\lambda}^0}}\sim\mathcal{V}^{-\frac{1}{3}}, \tag{64}$$

where for $i=1,2$, $f_i=(e_1,e_2)$, (u_1,u_2). Incorporating the value of masses in Eq. (63), using $\tilde{f}\sim\mathcal{V}^{-\frac{23}{30}}$ as obtained in [37], and Calabi-Yau volume $\mathcal{V}\sim10^5$, the dominant contribution of the EDM of the electron will be given as

$$\frac{d_e}{e}\bigg|_{\lambda^0} \equiv \frac{\mathcal{V}^{\frac{2}{3}}m_{\frac{3}{2}}(\tilde{f}^2\mathcal{V}^{-\frac{8}{5}})}{(4\pi)^2}\times\tilde{f}\mathcal{V}^{\frac{62}{45}}\left(\frac{\mathcal{V}^{-\frac{1}{3}}}{\mathcal{V}m^2_{\frac{3}{2}}}\right)$$
$$\equiv\frac{\tilde{f}^3\mathcal{V}^{\frac{2}{3}+\frac{62}{45}-\frac{8}{5}-\frac{1}{3}-1}}{(4\pi)^2m_{\frac{3}{2}}}$$
$$\equiv10^{-39}\ \text{cm}, \tag{65}$$

and the dominant contribution of the EDM of the neutron/quark will be given as

$$\frac{d_n}{e}\bigg|_{\lambda^0} \equiv \frac{\mathcal{V}^{\frac{2}{3}}m_{\frac{3}{2}}(\tilde{f}^2\mathcal{V}^{-\frac{7}{5}})}{(4\pi)^2}\times\tilde{f}\mathcal{V}^{\frac{62}{45}}\left(\frac{\mathcal{V}^{-\frac{1}{3}}}{\mathcal{V}m^2_{\frac{3}{2}}}\right)$$
$$\equiv\frac{\tilde{f}^3\mathcal{V}^{\frac{2}{3}+\frac{62}{45}-\frac{7}{5}-\frac{1}{3}-1}}{(4\pi)^2m_{\frac{3}{2}}}$$
$$\equiv10^{-38}\ \text{cm}. \tag{66}$$

2. Neutralino contribution

The physical eigenstates of the neutralino mass matrix in the context of the $\mathcal{N}=1$ gauged supergravity action are given as [38]

$$\tilde{\chi}^0_1\sim\frac{-\tilde{H}^0_u+\tilde{H}^0_d}{\sqrt{2}};\ m_{\chi^0_1}\sim\mathcal{V}^{\frac{59}{72}}m_{\frac{3}{2}},$$
$$\tilde{\chi}^0_2\sim\left(\frac{v}{M_P}\tilde{f}\mathcal{V}^{\frac{5}{6}}\right)\lambda^0+\frac{\tilde{H}^0_u+\tilde{H}^0_d}{\sqrt{2}};\ m_{\chi^0_2}\sim\mathcal{V}^{\frac{59}{72}}m_{\frac{3}{2}},$$
$$\tilde{\chi}^0_3\sim-\lambda^0+\left(\frac{v}{M_P}\tilde{f}\mathcal{V}^{\frac{5}{6}}\right)(\tilde{H}^0_u+\tilde{H}^0_d);\ m_{\chi^0_3}\sim\mathcal{V}^{\frac{2}{3}}m_{\frac{3}{2}}, \tag{67}$$

where v is value of the Higgs VEV at the electroweak scale. $\tilde{H}^0_u$ and $\tilde{H}^0_d$ correspond to an $SU(2)$-doublet Higgsino. $\tilde{\chi}^0_1$ is purely a Higgsino, and $\tilde{\chi}^0_2$ ($\tilde{\chi}^0_3$) are formed by a linear combination of a gaugino (Higgsino) with a very small admixture of Higgsino (gaugino). Since neutralinos are also very heavy, we evaluate the contribution of the same to the one-loop electron/neutron EDM involving heavy sfermions. Though the neutralino ($\chi^0_{1,2}$)-fermion-sfermion couplings are complex in this case, the phase disappears due to presence of both the complex coupling as well as its conjugate in the EDM expression. Therefore, the nonzero EDM arises due to CP-violating phases appearing from the mass eignstates of the sfermion mass matrix only. The one-loop diagram is given in Fig. 2.

We have already calculated the contribution of gaugino-lepton(quark)-slepton(quark) vertices in Sec. IV A. Now we estimate coefficients of the vertices corresponding to Higgsino-lepton(quark)-slepton(squark) interaction vertices.

[3] We only incorporate the volume suppression coming from $C_{\tilde{e}_i\tilde{e}_i^*\gamma}$ and $C_{\tilde{u}_i\tilde{u}_i^*\gamma}$. The momentum dependence of both vertices has already been included in the one-loop functions $A(r)$ and $B(r)$.

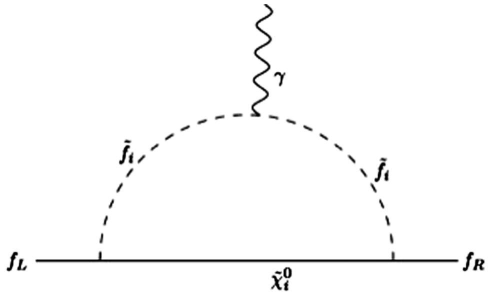

FIG. 2. One-loop diagram involving neutralinos.

In $\mathcal{N}=1$ gauged supergravity, Higgsino-fermion-sfermion interaction is governed by [39]

$$\mathcal{L}_{f-\tilde{f}-\tilde{H}^0_i} = \frac{e^{\frac{K}{2}}}{2}(\mathcal{D}_i D_{\bar{J}} W)\chi^i_L \chi^{j^c}_L + i g_{i\bar{J}}\bar{\chi}^{\bar{J}}\left[\bar{\sigma}\cdot\partial\chi^i + \Gamma^i_{Lk}\bar{\sigma}\cdot\partial a^L \chi^k + \frac{1}{4}(\partial_{a_L}K\bar{\sigma}\cdot a_L - \text{c.c.})\chi^i\right]. \quad (68)$$

Using this, we evaluate the coefficients of Higgsino-lepton(quark)-slepton(squark) interaction vertices. For an incoming electron (e^-_L) interacting with a slepton as well as a neutralino, the contribution of the Higgsino-lepton(e^-_L)-slepton($\tilde{e}_L$) vertex in the gauged supergravity action of Wess and Bagger [39] is given by $\frac{e^{\frac{K}{2}}}{2}(\mathcal{D}_{\mathcal{Z}_1}D_{\bar{\mathcal{A}}_1}W)\chi^{\mathcal{A}_i}\chi^{c^{\mathcal{Z}_I}} + i g_{\bar{I}\mathcal{A}_1}\bar{\chi}^{Z_i}[\bar{\sigma}\cdot\partial\chi^{\mathcal{A}_1} + \Gamma^{\mathcal{A}_1}_{\mathcal{A}_1\bar{A}_1}\bar{\sigma}\cdot\partial\mathcal{A}_1\chi^{\mathcal{A}_1} + \frac{1}{4}(\partial_{\mathcal{A}_3}K\bar{\sigma}\cdot\mathcal{A}_1 - \text{c.c.})\chi^{\mathcal{A}_1}]$. $\chi^{\mathcal{Z}}$ and $\chi^{c^{\mathcal{Z}_1}}$ correspond to an $SU(2)_L$ Higgsino and its charge conjugate, $\chi^{\mathcal{A}_1}$ corresponds to an $SU(2)_L$ electron, and $\tilde{\mathcal{A}}_1$ corresponds to the left-handed slepton. In the diagonalized set of basis, $g_{I\bar{\mathcal{A}}_1}=0$. Since $SU(2)$ EW symmetry is not conserved for the Higgsino-lepton-slepton vertex, to calculate the contribution of the same, we generate a term of the type $e_L\tilde{e}_L\tilde{H}^c_L H_L$ wherein e_L and H_L are, respectively, the $SU(2)_L$ electron and Higgs doublets, $\tilde{l}_L$ is also an $SU(2)_L$ doublet, and $\tilde{H}^c_L$ is an $SU(2)_L$ Higgsino doublet. After giving a VEV to one of the Higgs doublets H_L, one gets the required vertex. By considering $a_1 \to a_1 + \mathcal{V}^{-\frac{2}{9}}M_P$ and further picking up the component of $\mathcal{D}_i D_{\bar{a}_1}W$ linear in z_i as well as linear in fluctuation $(a_1 - \mathcal{V}^{-\frac{2}{9}}M_P)$, we see that $e^{\frac{K}{2}}\mathcal{D}_i D_{\bar{a}_1}W \sim \mathcal{V}^{-\frac{31}{18}}z_i(a_1 - \mathcal{V}^{-\frac{2}{9}}M_P)$. As was shown in [38], $e^{\frac{K}{2}}\mathcal{D}_I D_{\bar{\mathcal{A}}_1}W \sim \mathcal{O}(1)e^{\frac{K}{2}}\mathcal{D}_i D_{\bar{a}_1}W$. Utilizing the same, the magnitude of the physical Higgsino($\tilde{H}^c_L$)-lepton(e_L)-slepton($\tilde{e}_L$) vertex after giving a VEV to $\mathcal{Z}_I$ will be given as

$$|C_{\tilde{H}^c_L e_L \tilde{e}_L}| \equiv \frac{\mathcal{V}^{-\frac{31}{18}}\langle\mathcal{Z}_I\rangle}{\sqrt{\hat{K}^2_{\mathcal{Z}_1\bar{\mathcal{Z}}_1}\hat{K}_{\mathcal{A}_1\bar{\mathcal{A}}_1}\hat{K}_{\mathcal{A}_1\bar{\mathcal{A}}_1}}}\tilde{\mathcal{A}}_1\chi^{c^{\mathcal{Z}_I}}\chi^{\mathcal{A}_1} \equiv \mathcal{V}^{-\frac{3}{2}}. \quad (69)$$

The coefficient of the Higgsino($\tilde{H}^c_L$)-lepton(u_L)-slepton($\tilde{u_L}$) vertex can be determined by expanding $\mathcal{D}_I D_{\bar{\mathcal{A}}_2}W$ linear in $\mathcal{Z}_I$ as well as linear in fluctuation $(\mathcal{A}_2 - \mathcal{V}^{-\frac{1}{3}}M_P)$. The magnitude of the same has already been calculated in [38] and given as

$$|C_{\tilde{H}^c{}_L u_L \tilde{u_L}}| \equiv \frac{\mathcal{V}^{-\frac{20}{9}}\langle\mathcal{Z}_i\rangle}{\sqrt{\hat{K}^2_{\mathcal{Z}_1\bar{\mathcal{Z}}_1}\hat{K}_{\mathcal{A}_2\bar{\mathcal{A}}_2}\hat{K}_{\mathcal{A}_2\bar{\mathcal{A}}_2}}}\tilde{\mathcal{A}}_2\chi^{c^{\mathcal{Z}_I}}\chi^{\mathcal{A}_2} \equiv \mathcal{V}^{-\frac{4}{5}}. \quad (70)$$

To determine the contribution of the Higgsino-lepton(e^-_L)-slepton($\tilde{e}_R$) vertex, one needs to expand $\frac{e^{\frac{K}{2}}}{2}(\mathcal{D}_{\mathcal{Z}_I}D_{\bar{\mathcal{A}}_1}W)$ in the fluctuations linear in $\mathcal{A}_3$ about its stabilized value. Considering $a_3 \to a_3 + \mathcal{V}^{-\frac{13}{18}}M_P$ and picking up the component of $\mathcal{D}_I D_{\bar{a}_1}W$ linear in a_3, we have $e^{\frac{K}{2}}\mathcal{D}_i D_{\mathcal{A}_1}W \equiv e^{\frac{K}{2}}\mathcal{D}_I D_{\bar{a}_1}W \sim \mathcal{V}^{-\frac{43}{36}}(a_3 - \mathcal{V}^{-\frac{13}{18}}M_P)$. The contribution of the physical Higgsino($\tilde{H}^c_L$)-lepton(e_L)-slepton($\tilde{e_R}$) vertex will be given as

$$|C_{\tilde{H}^c_L e_L \tilde{e_R}}| \equiv \frac{\mathcal{V}^{-\frac{43}{36}}}{\sqrt{\hat{K}_{\mathcal{Z}_1\bar{\mathcal{Z}}_1}\hat{K}_{\mathcal{A}_1\bar{\mathcal{A}}_1}\hat{K}_{\mathcal{A}_3\bar{\mathcal{A}}_3}}}\tilde{\mathcal{A}}_3\chi^{Z_i}\chi^{c^{\mathcal{A}_1}} \equiv \mathcal{V}^{-\frac{9}{5}}. \quad (71)$$

Similarly, one can calculate the Higgsino($\tilde{H}^c_L$)-lepton(u_L)-slepton($\tilde{u_R}$) vertex by expanding $\frac{e^{\frac{K}{2}}}{2}(\mathcal{D}_{\mathcal{Z}_I}D_{\bar{\mathcal{A}}_2}W)$ in the fluctuations linear in $\mathcal{A}_4$ about its stabilized value. Considering $a_4 \to a_4 + \mathcal{V}^{-\frac{11}{9}}M_P$ and picking up the component of $\mathcal{D}_i D_{\bar{a}_1}W$ linear in a_4, we have $\frac{e^{\frac{K}{2}}}{2}(\mathcal{D}_{\mathcal{Z}_I}D_{\bar{\mathcal{A}}_2}W) \equiv e^{\frac{K}{2}}\mathcal{D}_i D_{\bar{a}_1}W \sim \mathcal{V}^{-\frac{43}{36}}(a_4 - \mathcal{V}^{-\frac{11}{9}}M_P)$. The coefficient of the Higgsino($\tilde{H}^c_L$)-lepton(u_L)-slepton($\tilde{u_R}$) vertex will be given as

$$|C_{\tilde{H}^c_L u_L \tilde{u_R}}| \equiv \frac{\mathcal{V}^{-\frac{43}{36}}}{\sqrt{\hat{K}_{\mathcal{Z}_1\bar{\mathcal{Z}}_1}\hat{K}_{\mathcal{A}_2\bar{\mathcal{A}}_2}\hat{K}_{\mathcal{A}_4\bar{\mathcal{A}}_4}}}\tilde{\mathcal{A}}_4\chi^{Z_i}\chi^{c^{\mathcal{A}_2}} \equiv \mathcal{V}^{-\frac{5}{3}}. \quad (72)$$

For an outgoing electron e^-_R interacting with a slepton as well as a neutralino, the contribution of the Higgsino-lepton(e^-_R)-slepton($\tilde{e}_L$) vertex is given by expanding $\frac{e^{\frac{K}{2}}}{2}(\mathcal{D}_{\mathcal{Z}_1}D_{\mathcal{A}_3}W)$ linear in $\mathcal{A}_1$. Considering $a_1 \to a_1 + \mathcal{V}^{-\frac{2}{9}}M_P$ and picking up the component of $\mathcal{D}_i D_{a_3}W$ linear in a_1, we have $e^{\frac{K}{2}}\mathcal{D}_I D_{\mathcal{A}_3}W \equiv e^{\frac{K}{2}}\mathcal{D}_i D_{a_3}W \sim \mathcal{V}^{-\frac{43}{36}}(a_1 - \mathcal{V}^{-\frac{2}{9}}M_P)$. The contribution of the physical Higgsino($\tilde{H}^c_L$)-lepton(e_R)-slepton($\tilde{e_L}$) vertex will be given as

$$|C_{\tilde{H}^c_L e_R \tilde{e_L}}| \equiv \frac{\mathcal{V}^{-\frac{37}{72}}}{\sqrt{\hat{K}_{\mathcal{Z}_1\bar{\mathcal{Z}}_1}\hat{K}_{\mathcal{A}_1\bar{\mathcal{A}}_1}\hat{K}_{\mathcal{A}_3\bar{\mathcal{A}}_3}}}\tilde{\mathcal{A}}_1\chi^{Z_i}\chi^{c^{\mathcal{A}_3}} \equiv \mathcal{V}^{-\frac{9}{5}}. \quad (73)$$

Similarly, considering $a_4 \to a_4 + \mathcal{V}^{-\frac{11}{9}} M_P$ and picking up the component of above term linear in a_1, we have $e^{\frac{K}{2}} \mathcal{D}_I D_{\mathcal{A}_4} W \equiv e^{\frac{K}{2}} \mathcal{D}_i D_{a_4} W \sim \mathcal{V}^{-\frac{43}{36}} (a_2 - \mathcal{V}^{-\frac{1}{3}} M_P)$. The contribution of physical Higgsino($\tilde{H}_L$)-quark(u_R)-squark($\tilde{u_L}$) vertex is given as

$$|C_{\tilde{H}_L u_R \tilde{u_L}}| \equiv \frac{\mathcal{V}^{-\frac{43}{36}}}{\sqrt{\hat{K}_{\mathcal{Z}_1 \bar{\mathcal{Z}}_1} \hat{K}_{\mathcal{A}_2 \bar{\mathcal{A}}_2} \hat{K}_{\mathcal{A}_4 \bar{\mathcal{A}}_4}}} \tilde{A}_2 \chi^{Z_i} \chi^{c^{\mathcal{A}_4}} \equiv \mathcal{V}^{-\frac{5}{3}}. \quad (74)$$

The Higgsino-lepton(e_R^-)-slepton($\tilde{e}_R$) vertex also does not possess $SU(2)$ EW symmetry. Therefore, to calculate the contribution of the same, we generate a term of the type $e_R \tilde{e}_R \tilde{H}_L H_L$, where H_L is one of the $SU(2)_L$ Higgs doublets. Thereafter, we expand $\frac{e^{\frac{K}{2}}}{2} (\mathcal{D}_{\mathcal{Z}_1} D_{\bar{\mathcal{A}}_3} W) \chi^{\mathcal{Z}_i} \chi^{\mathcal{A}_3}$ linear in $\mathcal{Z}_1$ and then linear in $\mathcal{A}_3$ about their stabilized VEVs. Considering $a_3 \to a_3 + \mathcal{V}^{-\frac{13}{18}} M_P$ and further picking up the component linear in z_i as well as linear in fluctuation $(a_3 - \mathcal{V}^{-\frac{2}{9}} M_P)$, we get $e^{\frac{K}{2}} \mathcal{D}_i D_{\mathcal{A}_3} W \equiv e^{\frac{K}{2}} \mathcal{D}_i D_{\bar{a}_3} W \sim \mathcal{V}^{-\frac{13}{18}} \langle z_i \rangle (a_3 - \mathcal{V}^{-\frac{13}{18}} M_P)$. The magnitude of the physical Higgsino($\tilde{H}_L$)-lepton(e_R)-slepton($\tilde{e_R}$) vertex after giving a VEV to $\mathcal{Z}_I$ is given as

$$|C_{\tilde{H}_L e_R \tilde{e_R}}| \equiv \frac{\mathcal{V}^{-\frac{13}{18}} \langle \mathcal{Z}_i \rangle}{\sqrt{\hat{K}^2_{\mathcal{Z}_1 \bar{\mathcal{Z}}_1} \hat{K}_{\mathcal{A}_3 \bar{\mathcal{A}}_3} \hat{K}_{\mathcal{A}_3 \bar{\mathcal{A}}_3}}} \tilde{\mathcal{A}}_3 \chi^{Z_I} \chi^{\mathcal{A}_3} \equiv \mathcal{V}^{-\frac{10}{9}}. \quad (75)$$

The contribution of the Higgsino-quark(u_R)-squark($\tilde{u}_R$) vertex has already been evaluated in [38] by expanding $\frac{e^{\frac{K}{2}}}{2} (\mathcal{D}_{\mathcal{Z}_I} D_{\mathcal{A}_4} W) \chi^{\mathcal{Z}_I} \chi^{\mathcal{A}_4}$ in the fluctuations linear in $\mathcal{Z}_I$ as well as $\mathcal{A}_4$ about their stabilized VEVs. The magnitude of the same is given as

$$|C_{\tilde{H}_L u_R \tilde{u_R}}| \equiv \frac{\mathcal{V}^{\frac{5}{18}} \langle \mathcal{Z}_i \rangle}{\sqrt{\hat{K}^2_{\mathcal{Z}_1 \bar{\mathcal{Z}}_1} \hat{K}_{\mathcal{A}_4 \bar{\mathcal{A}}_4} \hat{K}_{\mathcal{A}_4 \bar{\mathcal{A}}_4}}} \tilde{\mathcal{A}}_4 \chi^{Z_I} \chi^{\mathcal{A}_4} \equiv \mathcal{V}^{-\frac{10}{9}}. \quad (76)$$

The results of the coefficients of both slepton(squark)-lepton (quark)-Higgsino as given in Eqs. (69)–(76) are as follows:

$$\begin{aligned} &|C_{\tilde{H}^c_L e_L \tilde{e_L}}| \equiv \mathcal{V}^{-\frac{3}{2}}, \qquad |C_{\tilde{H}^c_L e_L \tilde{e_R}}| \equiv |C_{\tilde{H}_L e_R \tilde{e_L}}| \equiv \mathcal{V}^{-\frac{9}{5}}, \qquad |C_{\tilde{H}_L e_R \tilde{e_R}}| \equiv \mathcal{V}^{-\frac{10}{9}}, \\ &|C_{\tilde{H}^c_L u_L \tilde{u_L}}| \equiv \mathcal{V}^{-\frac{4}{5}}, \qquad |C_{\tilde{H}^c_L u_L \tilde{u_R}}| \equiv |C_{\tilde{H}_L u_R \tilde{u_L}}| \equiv \mathcal{V}^{-\frac{5}{3}}, \qquad |C_{\tilde{H}_L u_R | \tilde{u_R}}| \equiv \mathcal{V}^{-\frac{10}{9}}. \end{aligned} \quad (77)$$

Utilizing the aforementioned results and the results of various gaugino-fermion-sfermion vertices as given in Eq. (58), and by adding the contribution of the same according to Eq. (67), the volume-suppression factors coming from the neutralino-lepton-slepton vertices are given as

$$\begin{aligned} &|C_{\chi^0_1 e_L \tilde{e_L}}| = |C_{\chi^0_2 e_L \tilde{e_L}}| \equiv \mathcal{V}^{-\frac{3}{2}}, \qquad |C_{\chi^0_3 e_L \tilde{e_L}}| \equiv \tilde{f} \mathcal{V}^{-1}, \qquad |C_{\chi^0_1 e_L \tilde{e_R}}| = |C_{\chi^0_2 e_L \tilde{e_R}}| \equiv \mathcal{V}^{-\frac{9}{5}}, \\ &|C_{\chi^0_3 e_L \tilde{e_R}}| \equiv \tilde{f} \mathcal{V}^{-\frac{15}{9}}, \qquad |C_{\chi^0_1 e_R \tilde{e_L}}| = |C_{\chi^0_2 e_R \tilde{e_L}}| \equiv \mathcal{V}^{-\frac{9}{5}}, \qquad |C_{\chi^0_3 e_R \tilde{e_L}}| \equiv \tilde{f} \mathcal{V}^{-\frac{15}{9}}, \\ &|C_{\chi^0_1 e_R \tilde{e_R}}| = |C_{\chi^0_2 e_R \tilde{e_R}}| \equiv \mathcal{V}^{-\frac{10}{9}}, \qquad |C_{\chi^0_3 e_R \tilde{e_R}}| \equiv \tilde{f} \mathcal{V}^{-\frac{3}{5}}. \end{aligned} \quad (78)$$

The volume-suppression factors coming from the neutralino-quark-squark vertices are given as

$$\begin{aligned} &|C_{\chi^0_1 u_L \tilde{u_L}}| = |C_{\chi^0_2 u_L \tilde{u_L}}| \equiv \mathcal{V}^{-\frac{4}{5}}, \qquad |C_{\chi^0_3 u_L \tilde{u_L}}| \equiv \tilde{f} \mathcal{V}^{-\frac{4}{5}}, \qquad |C_{\chi^0_1 u_L \tilde{u_R}}| = |C_{\chi^0_2 u_L \tilde{u_R}}| \equiv \mathcal{V}^{-\frac{5}{3}}, \\ &|C_{\chi^0_3 u_L \tilde{u_R}}| \equiv \tilde{f} \mathcal{V}^{-\frac{14}{9}} |C_{\chi^0_1 u_R \tilde{u_L}}| = |C_{\chi^0_2 u_R \tilde{u_L}}| \equiv \mathcal{V}^{-\frac{5}{3}}, \qquad |C_{\chi^0_3 u_R \tilde{u_L}}| \equiv \tilde{f} \mathcal{V}^{-\frac{14}{9}}, \\ &|C_{\chi^0_1 u_R \tilde{u_R}}| = |C_{\chi^0_2 u_R \tilde{u_R}}| \equiv \mathcal{V}^{-\frac{10}{9}}, \qquad |C_{\chi^0_3 u_R \tilde{u_R}}| \equiv \tilde{f} \mathcal{V}^{-\frac{3}{5}}. \end{aligned} \quad (79)$$

The interaction Lagrangian governing the neutralino-slepton(squark)-lepton(quark) interaction can be written as

$$\mathcal{L} = \sum_{i=1,3} C_{\chi^0_i f_L \tilde{f_L}} f_L \tilde{f}_L \chi^0_i + C_{\chi^0_i f_L \tilde{f_R}} f_L \tilde{f}_R \chi^0_i + C_{\chi^0_i f_R \tilde{f_L}} f_R \tilde{f}_L \chi^0_i + C_{\chi^0_i f_R \tilde{f_R}} f_R \tilde{f}_R \chi^0_i, \quad (80)$$

where $f = (e, u)$. Rewriting f_L as well as f_R in term of the diagonalized basis states f_1 and f_2, the equation takes the form of Eq. (35):

$$\begin{aligned} \mathcal{L}_{\rm int} = \bar{\chi}_f \left((C_{\chi^0_i f_L \tilde{f_L}} D_{f_{11}} + C_{\chi^0_i f_L \tilde{f_R}} D_{f_{21}}) \frac{1+\gamma_5}{2} + (C_{\chi^0_i f_R \tilde{f_L}} D_{f_{11}} + C_{\chi^0_i f_R \tilde{f_R}} D_{f_{21}}) \frac{1-\gamma_5}{2} \right) \phi_{f_1} \chi^0_i \\ + \bar{\chi}_f \left((C_{\chi^0_i f_L \tilde{f_L}} D_{f_{12}} + C_{\chi^0_i f_L \tilde{f_R}} D_{f_{22}}) \frac{1+\gamma_5}{2} + (C_{\chi^0_i f_R \tilde{f_L}} D_{f_{12}} + C_{\chi^0_i f_R \tilde{f_R}} D_{f_{22}}) \frac{1-\gamma_5}{2} \right) \phi_{f_2} \chi^0_i. \end{aligned} \quad (81)$$

Using Eq. (36), the dipole moment contribution will follow:

$$\left.\frac{d_f}{e}\right|_{\chi_i} = \sum_{i=1,3}\frac{m_{\tilde{\chi}_i^0}}{(4\pi)^2}\left[\frac{1}{m^2_{\tilde{f}_1}}\mathrm{Im}(C_{\chi_i^0 f_L\tilde{f}_L}C_{\chi_i^0 f_R\tilde{f}_R}D_{f_{11}}D^*_{f_{21}} + C_{\chi_i^0 f_L\tilde{f}_R}C_{\chi_i^0 f_R\tilde{f}_L}D_{f_{21}}D^*_{f_{11}})Q'_{\tilde{f}_1}B\left(\frac{m^2_{\tilde{\chi}_i^0}}{m_{\tilde{f}_1}{}^2}\right)\right.$$
$$\left. + \frac{1}{m^2_{\tilde{f}_2}}\mathrm{Im}(C_{\chi_i^0 f_L\tilde{f}_L}C_{\chi_i^0 f_R\tilde{f}_R}D_{f_{12}}D^*_{f_{22}} + C_{\chi_i^0 f_L\tilde{f}_R}C_{\chi_i^0 f_R\tilde{f}_L}D_{f_{22}}D^*_{f_{12}})Q'_{\tilde{f}_2}B\left(\frac{m^2_{\tilde{\chi}_i^0}}{m_{\tilde{f}_2}{}^2}\right)\right]. \tag{82}$$

Using the values of the first-generation scalar/slepton mass $m_{\tilde{f}_1} = \mathcal{V}^{\frac{1}{2}}m_{\frac{3}{2}}$ and $m_{\tilde{\chi}_1^0} = m_{\tilde{\chi}_2^0} = \mathcal{V}^{\frac{59}{72}}m_{\frac{3}{2}}$, $m_{\tilde{\chi}_3^0} = \mathcal{V}^{\frac{2}{3}}m_{\frac{3}{2}}$, one gets

$$B\left(\frac{m^2_{\tilde{\chi}_2^0}}{m_{\tilde{f}_i}{}^2}\right) = B\left(\frac{m^2_{\tilde{\chi}_1^0}}{m_{\tilde{f}_i}{}^2}\right) = \frac{1}{2(\frac{m^2_{\tilde{\chi}_1^0}}{m_{\tilde{f}_i}{}^2}-1)^2}\left(1+\frac{m^2_{\tilde{\chi}_1^0}}{m_{\tilde{f}_i}{}^2}+\frac{m^2_{\tilde{\chi}_1^0}}{m_{\tilde{f}_i}{}^2}\ln\left(\frac{m^2_{\tilde{\chi}_1^0}}{m_{\tilde{f}_i}{}^2}\right)\right)\left(1-\frac{m^2_{\tilde{\chi}_1^0}}{m_{\tilde{f}_i}{}^2}\right) \sim \frac{1}{\mathcal{V}^{\frac{23}{36}}}, \tag{83}$$

$$B\left(\frac{m^2_{\tilde{\chi}_3^0}}{m_{\tilde{f}_i}{}^2}\right) = \frac{1}{2\left(\frac{m^2_{\tilde{\chi}_3^0}}{m_{\tilde{f}_i}{}^2}-1\right)^2}\left(1+\frac{m^2_{\tilde{\chi}_3^0}}{m_{\tilde{f}_i}{}^2}+\frac{m^2_{\tilde{\chi}_3^0}}{m_{\tilde{f}_i}{}^2}\ln\left(\frac{m^2_{\tilde{\chi}_3^0}}{m_{\tilde{f}_i}{}^2}\right)\right)\left(1-\frac{m^2_{\tilde{\chi}_3^0}}{m_{\tilde{f}_i}{}^2}\right) \sim \frac{m_{\tilde{f}_i}{}^2}{m^2_{\tilde{\chi}_3^0}} = \mathcal{V}^{-\frac{1}{3}}. \tag{84}$$

Utilizing above and the results of $C_{\chi_i^0 e_{L/R}\tilde{e_{L/R}}}$ as given in Eq. (78) and further simplifying, the dominant contribution of the EDM of the electron will be given as[4]

$$\left.\frac{d_e}{e}\right|_{\chi_i} \equiv \frac{\mathcal{V}^{\frac{59}{72}}m_{\frac{3}{2}}(\mathcal{V}^{-\frac{8}{3}}\sin\theta_e\sin\phi_e)}{(4\pi)^2\mathcal{V}^{\frac{23}{36}}}\left[\frac{C_{\tilde{e_2}\tilde{e_2}^*\gamma}}{m^2_{\tilde{e}_2}} - \frac{C_{\tilde{e_1}\tilde{e_1}^*\gamma}}{m^2_{\tilde{e}_1}}\right]. \tag{85}$$

Similarly, using the results of $C^{\chi_i^0 u_{L/R}\tilde{u_{L/R}}}$ as given in Eq. (79), the dominant contribution of the EDM of the quark will be given as

$$\left.\frac{d_u}{e}\right|_{\chi_i} \equiv \frac{\mathcal{V}^{\frac{59}{72}}m_{\frac{3}{2}}(\mathcal{V}^{-\frac{17}{9}}\sin\theta_e\sin\phi_e)}{(4\pi)^2\mathcal{V}^{\frac{23}{36}}}\left[\frac{C_{\tilde{u_2}\tilde{u_2}^*\gamma}}{m^2_{\tilde{u}_2}} - \frac{C_{\tilde{u_1}\tilde{u_1}^*\gamma}}{m^2_{\tilde{u}_1}}\right]. \tag{86}$$

Incorporating the value of $C_{\tilde{e}_i\tilde{e}_i\gamma}$ from Eq. (56), one gets

$$\left.\frac{d_e}{e}\right|_{\chi_i} \equiv \frac{\mathcal{V}^{\frac{59}{72}}m_{\frac{3}{2}}(\mathcal{V}^{-\frac{8}{3}}\sin\theta_e\sin\phi_e)}{(4\pi)^2\mathcal{V}^{\frac{23}{36}}}\mathcal{V}^{\frac{62}{45}}\tilde{f}\left[\frac{\cos^2\theta_e}{m^2_{\tilde{e}_2}} - \frac{\sin^2\theta_e}{m^2_{\tilde{e}_1}}\right],$$
$$\left.\frac{d_u}{e}\right|_{\chi_i} \equiv \frac{\mathcal{V}^{\frac{59}{72}}m_{\frac{3}{2}}(\mathcal{V}^{-\frac{17}{9}}\sin\theta_u\sin\phi_u)}{(4\pi)^2\mathcal{V}^{\frac{23}{36}}}\mathcal{V}^{\frac{62}{45}}\tilde{f}\left[\frac{\cos^2\theta_u}{m^2_{\tilde{u}_2}} - \frac{\sin^2\theta_u}{m^2_{\tilde{u}_1}}\right]. \tag{87}$$

Incorporating the value of $\sin\theta_e = \sin\theta_u = 1$, $\sin\phi_e = \sin\phi_u = (0,1]$, $\tilde{f} \sim \mathcal{V}^{-\frac{23}{30}}$, and value of scalar masses $m_{\tilde{e}_i} = m_{\tilde{u}_i} = \mathcal{V}^{\frac{1}{2}}m_{\frac{3}{2}}$, the numerical value of the EDM of the electron for this case will be

$$\left.\frac{d_e}{e}\right|_{\chi_i} \equiv \frac{\mathcal{V}^{\frac{59}{72}}m_{\frac{3}{2}}}{(4\pi)^2\mathcal{V}^{\frac{23}{36}}}\left(\tilde{f}\mathcal{V}^{\frac{62}{45}}\right)\times\mathcal{V}^{-\frac{8}{3}}\left(\frac{1}{\mathcal{V}m^2_{\frac{3}{2}}}\right)$$
$$\equiv \frac{\tilde{f}\mathcal{V}^{\frac{59}{72}+\frac{62}{45}-\frac{8}{3}-\frac{23}{36}-1}}{(4\pi)^2 m_{\frac{3}{2}}}$$
$$\equiv 10^{-37}\ \mathrm{cm}, \tag{88}$$

and the numerical value of the EDM of the neutron/quark will be

$$\left.\frac{d_n}{e}\right|_{\chi_i} \equiv \frac{\mathcal{V}^{\frac{59}{72}}m_{\frac{3}{2}}}{(4\pi)^2\mathcal{V}^{\frac{23}{36}}}\left(\tilde{f}\mathcal{V}^{\frac{62}{45}}\right)\times\mathcal{V}^{-\frac{17}{9}}\left(\frac{1}{\mathcal{V}m^2_{\frac{3}{2}}}\right)$$
$$\sim \frac{\tilde{f}\mathcal{V}^{\frac{59}{72}+\frac{62}{45}-\frac{17}{9}-\frac{23}{36}-1}}{(4\pi)^2 m_{\frac{3}{2}}}$$
$$\equiv 10^{-34}\ \mathrm{cm}. \tag{89}$$

3. R-parity violating vertices contribution

We have explicitly taken into account the contribution of R-parity violating couplings in the context of the $\mathcal{N}=1$ gauged supergravity limit of μ-split SUSY setup discussed in [38]. Although one would certainly expect a very suppressed value of the EDM because of the presence of heavy sfermions as well as vanishing contribution of R-parity violating vertices, we discuss the effect of the same on the EDM of the electron/neutron just to compare the order of magnitude of the EDM with respect to the R-parity conserving loop diagrams. Though the R-parity violating interaction vertices are complex but due to the presence of both a R-parity violating vertex as well as its

[4] We use the assumption that the complex phases appearing in the effective Yukawa couplings are of $\mathcal{O}(1)$.

conjugate in the one-loop diagrams as given in Fig. 3, the complex phase disappears and, therefore, contribution of the same to the EDM will vanish. However, similar to the neutralino and gaugino one-loop diagrams, the nonzero phase corresponding to the CP-violating effect can appear only by considering the chirality flip between the slepton (squark) fields appearing as propagators in the one-loop diagram. Because of chirality flip, the matrix amplitude depends on the off-diagonal component of the slepton (squark) mass matrix, the contribution of which further depends on complex trilinear coupling $\mathcal{A}_{IJK}$ as well as supersymmetric mass parameter μ.

The one-loop Feynman diagrams for the electron EDM mediated by R-parity violating interaction vertices are given in Fig. 3. Using the analytical results as given in [63] to get the numerical estimate of the EDM of the electron, we have

$$\begin{aligned}\frac{d_e}{e}\bigg|_{\rm RPV} =& -|C_{e_L\tilde{u}^c_R d_L}|^2 C_{\tilde{u}^c_R\tilde{u}^{c*}_R\gamma}\frac{2e}{3}|\mathcal{A}_{u_j}|\frac{m_{d_k}}{m^3_{\tilde{u}}}\sin\theta_u\sin\phi_{A_u}B(r_{d_k})\\ &-|C_{e_L\tilde{d}_L u^c_R}|^2 C_{\tilde{d}_L{}^*\tilde{d}_L\gamma}\frac{e}{3}|\mathcal{A}_{d_j}|\frac{m_{u_k}}{m^3_{\tilde{d}}}\sin\theta_d\sin\phi_{A_d}B(r_{u_k})\\ &-|C_{e_L\tilde{u}^c_R d_L}|^2 C_{u^c_R u^{c*}_R\gamma}\frac{2e}{3}|\mathcal{A}_{d_j}|\frac{m_{u_j}}{m^3_{\tilde{d}}}\sin\theta_d\sin\phi_{A_d}A(r_{u_k})\\ &-|C_{e_L\tilde{d}_L u^c_R}|^2 C_{d^*_L d_L\gamma}\frac{e}{3}|\mathcal{A}_{u_j}|\frac{m_{d_j}}{m^3_{\tilde{u}}}\sin\theta_u\sin\phi_{A_u}A(r_{d_k}),\end{aligned} \tag{90}$$

where $r_{(u_k/d_k)} = m^2_{(u_k/d_k)}/m^2_{\tilde{f}}$ and the form of one-loop functions $A(r)$ and $B(r)$ is defined in (33).

One can draw the similar R-parity violating one-loop diagram to calculate the quark EDM (u) by replacing $(e_L, e_R) \leftrightarrow (u_L, u_R)$ and $\tilde{u}^c \leftrightarrow \tilde{e}^c$. The analytical expression in the case of the quark EDM will be of the following form:

$$\begin{aligned}\frac{d_u}{e}|_{\rm RPV} =& -|C_{u_L\tilde{e}^c_R d_L}|^2 C_{\tilde{e}^c_R\tilde{e}^{c*}_R\gamma}\,e\,|\mathcal{A}_{e_j}|\frac{m_{d_k}}{m^3_{\tilde{e}}}\sin\theta_e\sin\phi_{A_e}B(r_{d_k})\\ &-|C_{u_L\tilde{d}_L e^c_R}|^2 C_{\tilde{d}^*_L\tilde{d}_L\gamma}\frac{e}{3}|\mathcal{A}_{d_j}|\frac{m_{e_k}}{m^3_{\tilde{d}}}\sin\theta_d\sin\phi_{A_d}B(r_{e_k})\\ &-|C_{u_L\tilde{e}^c_R d_L}|^2 C_{e^c_R e^{c*}_R\gamma}\frac{e}{3}|\mathcal{A}_{d_j}|\frac{m_{e_j}}{m^3_{\tilde{d}}}\sin\theta_d\sin\phi_{A_d}A(r_{e_k})\\ &-|C_{u_L\tilde{d}_L e^c_R}|^2 C_{d^*_L d_L\gamma}\,e\,|\mathcal{A}_{e_j}|\frac{m_{d_j}}{m^3_{\tilde{e}}}\sin\theta_e\sin\phi_{A_u}A(r_{d_k}).\end{aligned} \tag{91}$$

The magnitude of the coefficient of interaction vertices $C_{e_L\tilde{u}^c d}$, $C_{e_L\tilde{d}u^c}$, $C_{u_L\tilde{e}^c_R d_L}$, and $C_{u_L\tilde{d}_L e^c_R}$ have already been obtained in [38] and given as

$$C_{e_L\tilde{u}^c_R d_L} = C_{e_L\tilde{d}u^c_R} = C_{u_L\tilde{e}^c_R d_L} = C_{u_L\tilde{d}_L e^c_R} \equiv \mathcal{V}^{\frac{5}{3}} e^{i\phi_{y_\alpha}}, \tag{92}$$

where ϕ_{y_α} is the phase factor associated with complex R-parity violating interaction vertices.

The volume-suppression factors coming from the $C_{\tilde{u}^c_R\tilde{u}^{c*}_R\gamma}$, $C_{\tilde{e}^c_R\tilde{e}^{c*}_R\gamma}$, and $C_{\tilde{d}_L\tilde{d}_L{}^*\gamma}$ vertices have already been obtained in the case of gaugino one-loop diagrams and given as

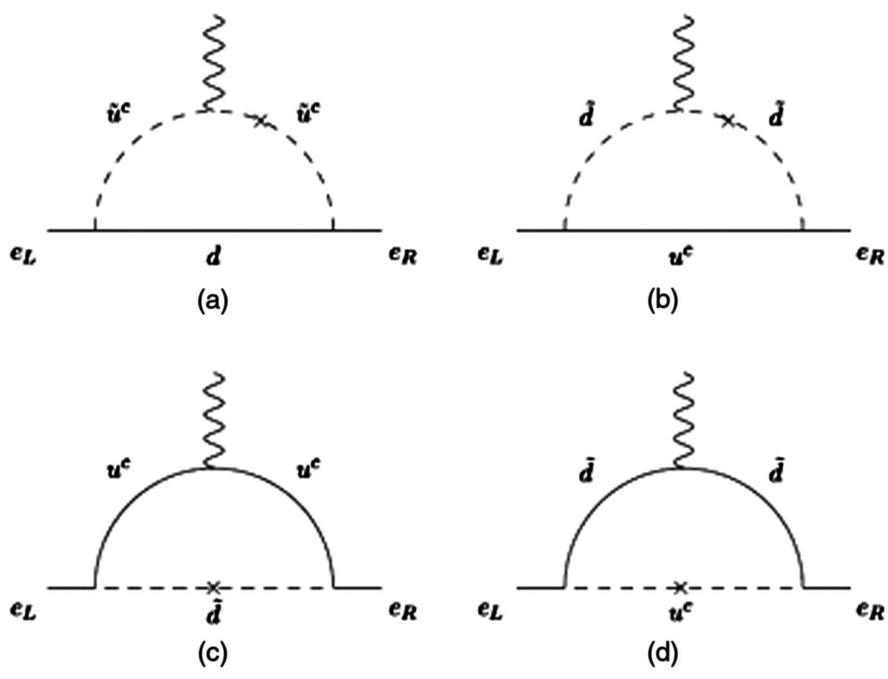

FIG. 3. One-loop diagrams involving R-parity violating couplings.

$$C_{\tilde{u}^c_R\tilde{u}^{c*}_R\gamma} = C_{\tilde{e}^c_R\tilde{e}^{c*}_R\gamma} \equiv \mathcal{V}^{\frac{62}{45}}\tilde{f}, \qquad C_{\tilde{d}_L\tilde{d}_L{}^*\gamma} \equiv \mathcal{V}^{\frac{53}{45}}\tilde{f}. \tag{93}$$

We set $C_{ff^*\gamma}|_{\rm EW}$ to be the charge of the quark d_L. The reason for the same is as follows. Consider the following kinetic-term-like term contributing to the quark-quark-photon vertex in the $\mathcal{N}=1$ gauged supergravity action of Wess and Bagger: $g_{YM}g_{a_2\bar{a}_2}\frac{\bar{\chi}_L^{\bar{a}_2}\Gamma^{a_2}_{T_Ba_2}X^B\not{A}\chi^{a_2}}{(\sqrt{K_{a_2\bar{a}_2}})^2} \in \frac{g_{a_2\bar{a}_2}\bar{\chi}^{\bar{a}_2}\not{D}\chi^{a_2}}{(\sqrt{K_{a_2\bar{a}_2}})^2}$. For the purpose of demonstrating the possibility of obtaining a SM-like quark-quark-photon coupling at the EW scale, let us assume that all moduli save T_B, a_2 have been stabilized at values indicated earlier and $(n^0_\beta)_{\rm max} \sim \mathcal{V}$ and, consequently, we take the Kähler potential to be

$$K \sim -2\ln\left[\left(T_B + \bar{T}_B - a_2\{C_{2\bar{2}}\bar{a}_2 + C_{2\bar{1}}\langle\bar{a}_1\rangle + C_{2\bar{3}}\langle\bar{a}_3\rangle + C_{2\bar{4}}\langle\bar{a}_4\rangle\} + {\rm c.c.} + \mathcal{V}^{\frac{2}{3}}\right)^{\frac{3}{2}} + \mathcal{V}\right]$$
$$\equiv -2\ln\left[\left(T_B + \bar{T}_B - C_{2\bar{2}}|a_2|^2 - a_2\bar{\Sigma}_2 + {\rm H.c.} + \mathcal{V}^{\frac{2}{3}}\right)^{\frac{3}{2}} + \mathcal{V}\right]. \tag{94}$$

Consider having frozen all moduli save T_B and a_2. Then from

$$\begin{pmatrix} g_{T_B\bar{T}_B} & g_{T_B\bar{a}_2} \\ g_{a_2\bar{T}_B} & g_{a_2\bar{a}_2}\end{pmatrix}^{-1} = \frac{1}{g_{T_B\bar{T}_B}g_{a_2\bar{a}_2} - |g_{T_B\bar{a}_2}|^2} \times \begin{pmatrix} g_{a_2\bar{a}_2} & -g_{T_B\bar{a}_2} \\ -g_{a_2\bar{T}_B} & g_{T_B\bar{T}_B}\end{pmatrix},$$

if $g_{T_B\bar{a}_2}|_{\rm EW}$ is small such that

$$|g_{T_B\bar{a}_2}|^2_{\rm EW} > g_{T_B\bar{T}_B}g_{a_2\bar{a}_2}|_{\rm EW,} \tag{95}$$

then

$$g^{T_B\bar{T}_B}|_{\rm EW} \sim \frac{g_{a_2\bar{a}_2}}{|g_{T_B\bar{a}_2}|^2}, \qquad g^{a_2\bar{a}_2}|_{\rm EW} \sim \frac{g_{T_B\bar{T}_B}}{|g_{T_B\bar{a}_2}|^2}|_{\rm EW},$$
$$g^{T_B\bar{a}_2}|_{\rm EW} \sim \frac{1}{g_{T_B\bar{a}_2}}|_{\rm EW} \equiv {\rm large}. \tag{96}$$

Using (94), we evaluate $\bar{\partial}_{\bar{a}_2}\partial_{a_2}K$, $\bar{\partial}_{\bar{T}_B}\partial_{T_B}K$, $\bar{\partial}_{\bar{T}_B}\partial_{a_2}K\partial_{T_B}\bar{\partial}$ and $\partial_{\bar{T}_B}\partial_{a_2}K$. If $\bar{\partial}_{\bar{T}_B}\partial_{a_2}K|_{\rm EW} \sim \delta \ll 1$ such that (95) is satisfied, then

$$3\langle(T_B + \bar{T}_B - C_{2\bar{2}}|a_2|^2 - a_2\bar{\Sigma}_2 + {\rm H.c.} + \mathcal{V}^{\frac{2}{3}})\rangle^{\frac{3}{2}}_{\rm EW}$$
$$\sim \langle[(T_B + \bar{T}_B - C_{2\bar{2}}|a_2|^2 - a_2\bar{\Sigma}_2 + {\rm H.c.} + \mathcal{V}^{\frac{2}{3}})^{\frac{3}{2}} + \mathcal{V}]\rangle_{\rm EW}. \tag{97}$$

Using (97), one sees that

$$\partial_{T_B}g_{a_2\bar{T}_B}|_{\rm EW\,near\,(97)} \sim \frac{9(C_{2\bar{2}}\bar{a}_2 + \bar{\Sigma}_2)}{[(T_B + \bar{T}_B - C_{2\bar{2}}|a_2|^2 - a_2\bar{\Sigma}_2 + {\rm H.c.} + \mathcal{V}^{\frac{2}{3}})^{\frac{3}{2}} + \mathcal{V}]^2} \sim \mathcal{V}^{-\frac{29}{18}}, \tag{98}$$

assuming $\langle a_{1,2,3,4}\rangle|_{\rm EW} \sim \mathcal{O}(1)\langle a_{1,2,3,4}\rangle|_{M_s}$. If $g_{a_2\bar{a}_2}|_{M_s} \sim g_{a_2\bar{a}_2}|_{\rm EW} \sim 10^{-2}$, then from (95), one sees

$$g_{T_B\bar{T}_B}|_{\rm EW} \sim \delta' < 10^2\delta^2. \tag{99}$$

Noting that

$$\Gamma^{a_2}_{T_ba_2} = \frac{g^{a_2\bar{T}_B}}{2}(\partial_{T_B}g_{a_2\bar{T}_B} + \partial_{a_2}g_{T_B\bar{T}_B}) + \frac{g^{a_2\bar{a}_2}}{2}(\partial_{T_B}g_{a_2\bar{a}_2} + \partial_{a_2}g_{T_B\bar{a}_2}), \tag{100}$$

we see that one can get a large contribution to (100) from $g^{a_2\bar{T}_B}\partial_{T_B}g_{a_2\bar{T}_B}|_{\rm EW}$ given by

$$g^{a_2\bar{T}_B}\partial_{T_B}g_{a_2\bar{T}_B}|_{{\rm EW},\mathcal{V}\sim 10^4} \sim \frac{10^{-6.5}}{\delta}. \tag{101}$$

Let us look at the implementation of (99) and its consequences. From the above calculations, one notes that (99) is identically satisfied if (97) is satisfied. Consider working with $\tau_{S,B}, z^i, a_I, \ldots$ instead of $T_{S,B}, z^i, a_I, \ldots$ having frozen G^a and other open-string moduli. Noting then that

$$K^{\alpha\bar{\beta}} = \frac{1}{K_{\tau_S\bar{\tau}_S}K_{\tau_B\bar{\tau}_B} - |K_{\tau_S\bar{\tau}_B}|^2}\begin{pmatrix} K_{\tau_B\bar{\tau}_B} & -K_{\tau_B\bar{\tau}_S} \\ -K_{\tau_S\bar{\tau}_B} & K_{\tau_S\bar{\tau}_S}\end{pmatrix}, \tag{102}$$

and assuming $K_{\tau_B\bar{\tau}_B}|_{\rm EW} \sim \delta' \ll 1$, $K_{\tau_S\bar{\tau}_S}|_{M_s} \sim K_{\tau_S\bar{\tau}_S}|_{\rm EW} \sim \mathcal{V}^{-1}$, $K_{\tau_S\bar{\tau}_B}|_{M_s} \sim K_{\tau_S\bar{\tau}_B}|_{\rm EW} \sim \mathcal{V}^{-\frac{5}{3}}$ implying $|K_{\tau_S\bar{\tau}_B}|^2 > K_{\tau_S\bar{\tau}_S}K_{\tau_B\bar{\tau}_B}$, one obtains

$$K^{\tau_S\bar{\tau}_S}|_{\rm EW} \sim \frac{K_{\tau_B\bar{\tau}_B}|_{\rm EW}}{|K_{\tau_S\bar{\tau}_B}|^2_{\rm EW}} \sim \delta'\mathcal{V}^{\frac{10}{3}}, \qquad K^{\tau_B\bar{\tau}_B}|_{\rm EW} \sim \frac{K_{\tau_S\bar{\tau}_S}|_{\rm EW}}{|K_{\tau_S\bar{\tau}_B}|^2_{\rm EW}} \sim \mathcal{V}^{\frac{7}{3}},$$
$$K^{\tau_S\bar{\tau}_B}|_{\rm EW} \sim \frac{1}{K_{\tau_S\bar{\tau}_B}|_{\rm EW}} \sim \mathcal{V}^{\frac{5}{3}}. \tag{103}$$

Equation (103) implies

$$\bar{F}^{\tau_S}|_{M_s} \sim K^{\tau_S\bar{\tau}_S} D_{\tau_S} W \sim \bar{F}^{\tau_S}|_{\rm EW} = e^{\frac{K}{2}}(K^{\tau_S\bar{\tau}_S} D_{\tau_S} W + K^{\tau_B\bar{\tau}_S} D_{\tau_B} W) \sim (\delta' \mathcal{V}^{\frac{10}{3}} + \mathcal{V}) m_{3/2} \sim \frac{1}{\mathcal{V}}. \quad (104)$$

So, the F^{τ_S} term (potential $\|F^{\tau_S}\|^2$) is one-loop RG invariant. Further, the complete F-term potential

$$V|_{M_s} \sim e^K K^{\tau_S\bar{\tau}_S} |D_{\tau_S} W|^2 \sim \mathcal{V} m_{3/2}^2 \sim e^K (K^{\tau_S\bar{\tau}_S}|D_{\tau_S}W|^2 + K^{\tau_B\bar{\tau}_B}|D_{\tau_B}W|^2 + K^{\tau_S\bar{\tau}_B} D_{\tau_S} D_{\bar{\tau}_B}\bar{W} + \text{H.c.})_{\rm EW} \sim (\delta' \mathcal{V}^{\frac{10}{3}} + \mathcal{V}) m_{3/2}^2 \sim \mathcal{V} m_{3/2}^2 \quad (105)$$

is also one-loop RG invariant. So, the quark-quark-photon vertex can be made to be of $\mathcal{O}(1)$ for $\delta \sim 10^{-13}$; i.e., one can hope that the coupling $C_{ff^*\gamma} \sim \mathcal{O}(1)$ for $f(\text{fermion}) \equiv e, u$.

For $r_{(u_k/d_k/e_k)} = m^2_{(u_k/d_k/e_k)}/m^2_{\tilde{f}_i}$, $A(r_{u_k/d_k/e_k}) = B(r_{u_k/d_k/e_k}) = 1$. As mentioned in Eq. (30), $|\mathcal{A}'_e| = |\mathcal{A}^*_e - \mu\cot\beta| \equiv \mathcal{V} m_{\frac{3}{2}}$, $|\mathcal{A}'_u| = |\mathcal{A}^*_u - \mu\cot\beta| \equiv \mathcal{V} m_{\frac{3}{2}}$. Using these results and the results of the coefficient of interaction vertices as given above, considering $\sin\phi_u = \sin\phi_d = (0,1]$, $\sin\theta_e = \sin\theta_u = 1$, the magnitude of the dominant contribution of the EDM of the electron will be given as

$$\frac{d_e}{e}\Big|_{\rm RPV} \sim \frac{2}{3}\frac{\mathcal{V}^{-\frac{10}{3}+1+\frac{62}{45}}}{\mathcal{V}^{\frac{3}{2}} m^2_{\frac{3}{2}}} m_{u_k} \equiv 10^{-31}\ \text{GeV}^{-1} \equiv 10^{-45}\ \text{cm}, \quad (106)$$

and the magnitude of dominant contribution of the EDM of the neutron/quark will be given as follows:

$$\frac{d_n}{e}\Big|_{\rm RPV} \sim \frac{\mathcal{V}^{-\frac{10}{3}+1+\frac{62}{45}}}{\mathcal{V}^{\frac{3}{2}} m^2_{\frac{3}{2}}} m_{e_k} \equiv 10^{-31}\ \text{GeV}^{-1} \equiv 10^{-45}\ \text{cm}. \quad (107)$$

B. One-loop diagrams involving neutral scalar (Higgs) in the loop

In this subsection, we estimate the contribution of one-loop diagrams involving fermions and Higgs as propagators to the EDM of a fermion. The fine-tuning argument given by Arkani-Hamed and Dimopoulos in [1] is not just able to provide a light Higgs by diagonalizing the Higgs mass matrix; it is important to give a reasonable order of magnitude of the EDM by considering diagonalized Higgs mass eigenstates (light Higgs as one of the eigenstates of the Higgs mass matrix) as scalar propagators in the one-loop diagram. In the discussion so far, we have argued that CP-violating phases in the one-loop diagram's contribution to the EDM of an electron/neutron are accomplished by considering the off-diagonal contribution of the sfermion mass matrix at the electroweak scale. In this subsection, we will discuss the one-loop diagrams in which nonzero CP-violating phases appear through mixing between the Higgs doublet in the Higgs mass matrix. Using the same approach, we have already calculated the mass of one of the Higgs formed by the linear combination of two Higgs doublets $H_{u,d}$ to be light (identified with position moduli $\mathcal{Z}_{1,2}$ in our setup; see [37,38]). Now, we implement this approach to calculate the nonzero EDM of the electron/neutron by considering the eigenstates of the Higgs mass matrix as propagators in the one-loop diagram.

1. SM-like Yukawa coupling contribution

The one-loop diagram mediated by a SM-like Yukawa coupling is given in Fig. 4. The effective one-loop operator given in Eq. (31) can be recasted in the following form:

$$\mathcal{L}_{\rm int} = \sum_i \bar{\chi}_f \left(C_{f_L^* f_R H_i}\frac{1-\gamma_5}{2} + C_{f_L^* f_R H_i}\frac{1+\gamma_5}{2}\right)\phi_{H_i}\chi_f + \text{H.c.} \quad (108)$$

For $i = 1, 2$, the above equation can be expanded as

$$\mathcal{L}_{\rm int} = \bar{\chi}_f \left(C_{f_L^* f_R H_1}\frac{1-\gamma_5}{2} + C_{f_L^* f_R H_1}\frac{1+\gamma_5}{2}\right)\phi_{H_1}\chi_f + \bar{\chi}_e \left(C_{f_L^* f_R H_2}\frac{1-\gamma_5}{2} + C_{f_L^* f_R H_2}\frac{1+\gamma_5}{2}\right)\phi_{H_2}\chi_f + \text{H.c.},$$

where ϕ_{H_1} and ϕ_{H_2} correspond to the eigenstates of the mass matrix of the Higgs doublet and χ_f corresponds to a fermion. Using Eq. (23), the aforementioned vertices can be expressed in terms of an undiagonalized (H_u, H_d) basis as follows:

$$C_{f_L^* f_R H_1} = D_{h_{11}} C_{f_L^* f_R H_u} + D_{h_{12}} C_{f_L^* f_R H_d},$$
$$C_{f_L^* f_R H_2} = D_{h_{21}} C_{f_L^* f_R H_u} + D_{h_{22}} C_{f_L^* f_R H_d}. \quad (109)$$

In $\mathcal{N} = 1$ gauged supergravity, the interaction vertices $C_{e_L^* e_R H_u/H_d}$ and $C_{u_L^* u_R H_u/H_d}$ will be given by expanding $e^{\frac{K}{2}} \mathcal{D}_{\mathcal{A}_1} D_{\mathcal{A}_3} W$ and $e^{\frac{K}{2}} \mathcal{D}_{\mathcal{A}_2} D_{\mathcal{A}_4} W$, respectively, in the

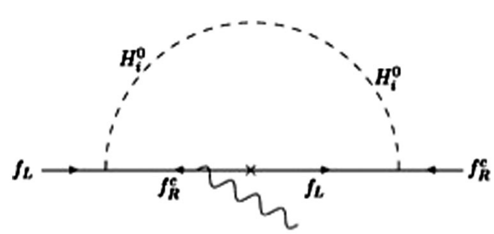

FIG. 4. One-loop diagram involving scalar (Higgs) and SM-like formions.

fluctuations linear in $\mathcal{Z}_i$ about its stabilized VEV. The values of the same have already been obtained in [38] and given as

$$\begin{aligned}
\text{for } f = e,\ C_{e_L^* e_R H_u/H_d} &= \hat{Y}^{\rm eff}_{\mathcal{Z}_I \mathcal{A}_1 \mathcal{A}_3} \\
&= \frac{\mathcal{O}(\mathcal{Z}_I - \mathcal{V}^{\frac{1}{36}}) \text{ term in } e^{\frac{K}{2}} \mathcal{D}_{\mathcal{A}_1} D_{\mathcal{A}_3} W}{\sqrt{K_{\mathcal{Z}_i \bar{\mathcal{Z}}_i} K_{\mathcal{A}_1 \bar{\mathcal{A}}_1} K_{\mathcal{A}_3 \bar{\mathcal{A}}_3}}} \\
&\equiv \mathcal{V}^{-\frac{47}{45}} e^{i\phi_{y_e}}, \\
\text{for } f = u,\ C_{u_L^* u_R H_u/H_d} &= \hat{Y}^{\rm eff}_{\mathcal{Z}_I \mathcal{A}_2 \mathcal{A}_4} \\
&= \frac{\mathcal{O}(\mathcal{Z}_I - \mathcal{V}^{\frac{1}{36}}) \text{ term in } e^{\frac{K}{2}} \mathcal{D}_{\mathcal{A}_2} D_{\mathcal{A}_4} W}{\sqrt{K_{\mathcal{Z}_1 \bar{\mathcal{Z}}_1} K_{\mathcal{A}_1 \bar{\mathcal{A}}_1} K_{\mathcal{A}_3 \bar{\mathcal{A}}_3}}} \\
&\equiv \mathcal{V}^{-\frac{19}{18}} e^{i\phi_{y_u}}, \qquad (110)
\end{aligned}$$

where $e^{i\phi_{y_e}}$ and $e^{i\phi_{y_e}}$ are the phase factors associated with complex effective Yukawa couplings.

Going back to Eq. (109),

$$\begin{aligned}
C_{e_L^* e_R H_1} &\equiv \mathcal{V}^{-\frac{47}{45}} e^{i\phi_{y_e}} (D_{h_{11}} + D_{h_{21}}), \\
C_{e_L^* e_R H_2} &\equiv \mathcal{V}^{-\frac{47}{45}} e^{i\phi_{y_e}} (D_{h_{12}} + D_{h_{22}}), \\
C_{u_L^* u_R H_1} &\equiv \mathcal{V}^{-\frac{19}{18}} e^{i\phi_{y_u}} (D_{h_{11}} + D_{h_{21}}), \\
C_{u_L^* u_R H_2} &\equiv \mathcal{V}^{-\frac{19}{18}} e^{i\phi_{y_u}} (D_{h_{12}} + D_{h_{22}}). \qquad (111)
\end{aligned}$$

Now, the one-loop EDM of the electron (quark) in this case will be given as [64]

$$\begin{aligned}
\frac{d}{e}\bigg|_{H_{1,2}} = \frac{m_f Q_f}{(4\pi)^2} \Bigg(&\frac{1}{m^2_{H_1}} \mathrm{Im}(C_{f_L^* f_R H_1} C^*_{f_L^* f_R H_1}) A\left(\frac{m_f^2}{m^2_{H_1}}\right) \\
&+ \frac{1}{m^2_{H_2}} \mathrm{Im}(C_{f_L^* f_R H_2} C^*_{f_L^* f_R H_2}) A\left(\frac{m_f^2}{m^2_{H_2}}\right)\Bigg), \qquad (112)
\end{aligned}$$

where m_f corresponds to the fermion mass, and $m_{H_{1,2}}$ corresponds to the eigenstates of the Higgs mass matrix. Since we are considering only first-generation fermions in our $D3/D7$ μ-split SUSY setup, the physical mass eigenstate of the fermion is the same as the usual Dirac mass term corresponding to the first-generation lepton/quark only. Using the fact that the phase factors associated with the Wilson line modulus $\mathcal{A}_{1/2}$ (identified with a first-generation L-hand lepton/quark), Wilson line modulus $\mathcal{A}_{3/4}$ (identified with a first-generation R-hand lepton/quark), and position modulus (identified with a Higgs doublet) are distinct, and the effective Yukawa couplings also produce a nonzero phase factor, the masses of SM fermions can be complex. Therefore, we assume that the overall phase formed by adding all phase factors associated with the fields and the coefficients of the Yukawa coupling add up in such a way that the overall phase vanishes and the fermion mass is real.

Using (111),

$$\begin{aligned}
\mathrm{Im}(C_{e_L^* e_R H_2} C^*_{e_L^* e_R H_2}) &= -\mathrm{Im}(C_{e_L^* e_R H_1} C^*_{e_L^* e_R H_1}) \\
&\equiv \frac{1}{2} \mathcal{V}^{-\frac{94}{45}} \sin\theta_h \sin\phi_h, \\
\mathrm{Im}(C_{u_L^* u_R H_2} C^*_{u_L^* u_R H_2}) &= -\mathrm{Im}(C_{u_L^* u_R H_1} C^*_{u_L^* u_R H_1}) \\
&\equiv \frac{1}{2} \mathcal{V}^{-\frac{19}{18}} \sin\theta_h \sin\phi_h.
\end{aligned}$$

Given that $\sin\theta_h = \frac{2|\hat{\mu}\mathcal{B}|}{\sqrt{(M^2_{H_u} - M^2_{H_d})^2 + 4(\hat{\mu}\mathcal{B})^2}}$, using the values given above, $\sin\theta_h \in [0,1]$. We also make an assumption that $\phi_h \ni (0, \frac{\pi}{2}]$. Using Eq. (33) and the value of $m_e = 0.5$ MeV, $m_{H_1} \equiv 125$ GeV, and $m_{H_2} \equiv \mathcal{V}^{\frac{59}{72}} m_{\frac{3}{2}}$, $A(\frac{m_e^2}{m^2_{H_1}}) = A(\frac{m_e^2}{m^2_{H_2}}) \equiv 1$. Using (33), the dominant contribution of the electron EDM in this case will be given as

$$\begin{aligned}
\frac{d_e}{e}\bigg|_{H_{1,2}} &= \frac{10^{-3}}{4(4\pi)^2} \mathcal{V}^{-\frac{94}{45}} \left(\frac{1}{m^2_{H_1}} - \frac{1}{m^2_{H_2}}\right) \\
&\equiv 10^{-20} \text{ GeV}^{-1} \\
&\equiv \mathcal{O}(10^{-34}) \text{ cm}. \qquad (113)
\end{aligned}$$

The numerical estimate of the neutron/quark EDM will be given as

$$\begin{aligned}
\frac{d_n}{e}\bigg|_{H_{1,2}} &= \frac{10^{-3}}{2(4\pi)^2} \mathcal{V}^{-\frac{19}{9}} \left(\frac{1}{m^2_{H_1}} - \frac{1}{m^2_{H_2}}\right) \\
&\equiv 10^{-29} \text{ GeV}^{-1} \\
&\equiv \mathcal{O}(10^{-33}) \text{ cm}. \qquad (114)
\end{aligned}$$

2. Chargino contribution

The one-loop diagram corresponding to the electron EDM mediated via Higgs and chargino exchange is given in Fig. 5. Because of the presence of heavy fermions and light as well as heavy scalars (eigenvalues of the Higgs mass matrix) existing as propagators in the loop, using an analytical expression of the one-loop EDM as given in Eq. (36), one can expect an enhancement in the order of magnitude of the EDM. We explicitly analyze the contribution of this loop diagram to the EDM at one loop in the

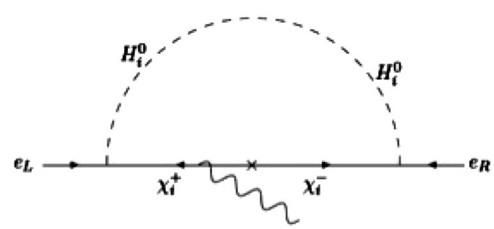

FIG. 5. One-loop diagram involving Higgs and charginos.

context of $\mathcal{N}=1$ gauged supergravity action. One cannot have a similar diagram for the quark because of the violation of charge conservation. So we use the loop diagram given in Fig. 5 to get the analysis of the EDM of the electron only. The effective one-loop operator will be of the following form:

$$\mathcal{L}_{\rm int} = \sum_{i,j}\bar{\chi}_f\left(C_{f_L^*\chi_j^+H_i^0}\frac{1-\gamma_5}{2}\right)\phi_{H_i^0}\tilde{\chi}_j^+ + \bar{\chi}_f\left(C_{q_R^*\chi_j^-H_i^0}\frac{1+\gamma_5}{2}\right)\phi_{H_i^0}\tilde{\chi}_j^- + \text{H.c.}\ldots \quad i,j=1,2. \tag{115}$$

Using Eq. (23), one can represent the coefficient of interaction vertices in terms of the undiagonalized basis of the Higgs mass matrix as follows:

$$\begin{aligned}
C_{e_L^*\chi_j^+H_1^0} &= D_{h_{11}}C_{e_L^*\chi_j^+H_u^0} + D_{h_{12}}C_{e_L^*\chi_j^+H_d^0},\\
C_{e_L^*\chi_j^+H_2^0} &= D_{h_{21}}C_{e_L^*\chi_j^+H_u^0} + D_{h_{22}}C_{e_L^*\chi_j^+H_u^0},\\
C_{e_R^*\chi_j^-H_1^0} &= D_{h_{11}}C_{e_R^*\chi_j^-H_u^0} + D_{h_{12}}C_{e_R^*\chi_j^-H_d^0},\\
C_{e_R^*\chi_j^-H_2^0} &= D_{h_{21}}C_{e_R^*\chi_j^-H_u^0} + D_{h_{22}}C_{e_R^*\chi_j^-H_u^0}.
\end{aligned} \tag{116}$$

As given in the Appendix,

$$\begin{aligned}
\tilde{\chi}_1^+ &= -\tilde{H}_u^+ + \left(\frac{v}{M_P}\tilde{f}\mathcal{V}^{\frac{5}{6}}\right)\tilde{\lambda}_i^+,\\
\tilde{\chi}_1^- &= -\tilde{H}_d^- + \left(\frac{v}{M_P}\tilde{f}\mathcal{V}^{\frac{5}{6}}\right)\tilde{\lambda}_i^-, \quad \text{and} \quad m_{\tilde{\chi}_1^\pm} \equiv \mathcal{V}^{\frac{59}{72}}m_{\frac{3}{2}},\\
\tilde{\chi}_2^+ &= \tilde{\lambda}_i^+ + \left(\frac{v}{M_P}\tilde{f}\mathcal{V}^{\frac{5}{6}}\right)\tilde{H}_u^+,\\
\tilde{\chi}_2^- &= \tilde{\lambda}_i^- + \left(\frac{v}{M_P}\tilde{f}\mathcal{V}^{\frac{5}{6}}\right)\tilde{H}_d^-, \quad \text{and} \quad m_{\tilde{\chi}_2^\pm} \equiv \mathcal{V}^{\frac{2}{3}}m_{\frac{3}{2}}.
\end{aligned}$$

Using the above,

$$\begin{aligned}
C_{e_L^*\chi_1^+H_u^0/H_d^0} &= -C_{e_L^*\tilde{H}_u^+H_u^0/H_d^0} + \left(\frac{v}{M_P}\tilde{f}\mathcal{V}^{\frac{5}{6}}\right)C_{e_L^*\tilde{\lambda}_i^+H_u^0/H_d^0},\\
C_{e_L^*\chi_2^+H_u^0/H_d^0} &= C_{e_L^*\tilde{\lambda}_i^+H_u^0/H_d^0} + \left(\frac{v}{M_P}\tilde{f}\mathcal{V}^{\frac{5}{6}}\right)C_{e_L^*\tilde{H}_u^+H_u^0/H_d^0},\\
C_{e_R^*\chi_1^-H_u^0/H_d^0} &= -C_{e_R^*\tilde{H}_d^-H_u^0/H_d^0} + \left(\frac{v}{M_P}\tilde{f}\mathcal{V}^{\frac{5}{6}}\right)C_{e_R^*\tilde{\lambda}_i^+H_u^0/H_d^0},\\
C_{e_R^*\chi_2^-H_u^0/H_d^0} &= C_{e_R^*\tilde{\lambda}_i^+H_u^0/H_d^0} + \left(\frac{v}{M_P}\tilde{f}\mathcal{V}^{\frac{5}{6}}\right)C_{e_R^*\tilde{H}_d^-H_u^0/H_d^0}.
\end{aligned} \tag{117}$$

The interaction vertices $C_{e_L^*\tilde{H}_u^+H_u^0/H_d^0}$ and $C_{e_R^*\tilde{H}_d^-H_u^0/H_d^0}$ corresponding to Fig. 5 will be given by expanding the $e^{\frac{K}{2}}\mathcal{D}_{\mathcal{Z}_1}D_{\mathcal{A}_1}W$ and $e^{\frac{K}{2}}\mathcal{D}_{\mathcal{Z}_1}D_{\mathcal{A}_3}W$ in the fluctuations linear in $\mathcal{Z}_i$ about its stabilized VEV. The contributions of $e^{\frac{K}{2}}\mathcal{D}_{z_1}D_{a_1}W$ as well as $e^{\frac{K}{2}}\mathcal{D}_{z_1}D_{a_3}W$ have been given in terms of the undiagonalized (z_i,a_i) basis in [38]. We assume that $e^{\frac{K}{2}}\mathcal{D}_iD_{\bar{\mathcal{A}}_1}W\sim\mathcal{O}(1)e^{\frac{K}{2}}\mathcal{D}_iD_{\bar{a}_1}W$. Since the EW symmetry gets broken for the Higgsino($\tilde{H}_u^+$)-lepton(e_L)-Higgs(H_u^0/H_d^0) vertex, we evaluate the contribution of the same by expanding $e^{\frac{K}{2}}\mathcal{D}_{z_1}D_{a_1}W$ in the fluctuations linear in z_1 as well as $(z_i-\mathcal{V}^{\frac{1}{36}})$ and then giving a VEV to z_i. Doing so, the magnitude of the coefficient of this vertex will be given as

$$|C_{e_L^*\tilde{H}_u^+H_u^0/H_d^0}| \sim \frac{\langle\mathcal{Z}_i\rangle\mathcal{O}(\mathcal{Z}_i-\mathcal{V}^{\frac{1}{36}})\ \text{term in}\ e^{\frac{K}{2}}\mathcal{D}_{\mathcal{Z}_1}D_{\mathcal{A}_1}W}{\sqrt{(K_{\mathcal{Z}_1\bar{\mathcal{Z}}_1})^3K_{\mathcal{A}_1\bar{\mathcal{A}}_1}}} \equiv \mathcal{V}^{-\frac{1}{10}}, \quad \text{for}\ \mathcal{V}=10^5. \tag{118}$$

Similarly, the contribution of the physical Higgsino($\tilde{H}_d^-$)-lepton(e_R)-Higgs(H_u^0/H_d^0) vertex will be given as

$$|C_{e_R^*\tilde{H}_d^-H_u^0/H_d^0}| \sim \frac{\mathcal{O}(\mathcal{Z}_i-\mathcal{V}^{\frac{1}{36}})\ \text{term in}\ e^{\frac{K}{2}}\mathcal{D}_{\mathcal{Z}_1}D_{\mathcal{A}_3}W}{\sqrt{K_{\mathcal{Z}_1\bar{\mathcal{Z}}_1}K_{\mathcal{Z}_1\bar{\mathcal{Z}}_1}K_{\mathcal{A}_3\bar{\mathcal{A}}_3}}} \equiv \mathcal{V}^{\frac{1}{10}}, \quad \text{for}\ \mathcal{V}=10^5. \tag{119}$$

The coefficient of the interaction vertex $e_L^- - H_u^0 - \tilde{\lambda}_i^+$ corresponding to Fig. 5 will be given by $\mathcal{L}_{e_L^- - H_u^0 - \tilde{\lambda}_i^+} = g_{YM}g_{\mathcal{A}_1\bar{T}_B}X^{*B}\bar{\chi}^{\bar{\mathcal{A}}_1}\tilde{\lambda}_i^+ + \partial_{\mathcal{A}_1}T_BD^B\bar{\chi}^{\bar{\mathcal{A}}_1}\tilde{\lambda}_i^+$. Since $\partial_{\mathcal{A}_1}T_B$ does not give any term which is linear in $\mathcal{Z}_i$, the second term contributes zero to the given vertex. By expanding $g_{\mathcal{A}_1\bar{T}_B}$ in the fluctuation linear in $\mathcal{Z}_1$ around its stabilized VEV, in terms of the undiagonalized basis, we have $g_{T_B\bar{a}_1}\to -\mathcal{V}^{-\frac{13}{12}}(z_1-\mathcal{V}^{-\frac{1}{36}})$ and $g_{YM}\sim\mathcal{V}^{-\frac{1}{36}}$. Considering $g_{YM}g_{T_B\bar{a}_1}\sim\mathcal{O}(1)g_{YM}g_{T_B\bar{\mathcal{A}}_1}$ as shown in [38], incorporating the values of $X^B=-6i\kappa_4^2\mu_7Q_{T_B}$, $\kappa_4^2\mu_7\sim\frac{1}{\mathcal{V}}$, and $Q_{T_B}\sim\mathcal{V}^{\frac{1}{3}}(2\pi\alpha')^2\tilde{f}$, we get the contribution of physical gaugino($\tilde{\lambda}_i^+$)-lepton(e_L)-Higgs(H_u^0) interaction vertex given as follows:

$$|C_{e_L\tilde{\lambda}_i^+H_u^0/H_d^0}| \equiv \frac{g_{YM}g_{T_B\bar{\mathcal{A}}_1}X^{T_B}\sim\mathcal{V}^{-\frac{47}{36}}\tilde{f}}{\sqrt{\hat{K}_{\mathcal{A}_1\bar{\mathcal{A}}_1}\hat{K}_{\mathcal{Z}_1\bar{\mathcal{Z}}_1}}}\mathcal{Z}_1\bar{\chi}^{\bar{\mathcal{A}}_1}\tilde{\lambda}^0 \equiv \tilde{f}(\mathcal{V}^{-\frac{3}{2}}). \tag{120}$$

To calculate the coefficient of interaction vertex $e_R^* - H_u^0 - \tilde{\lambda}_i^-$, we need to expand $g_{\mathcal{A}_3\bar{T}_B}$ in the fluctuation quadratic in $\mathcal{Z}_1$ to first conserve $SU(2)_L$ symmetry and after giving a VEV to one of the $\mathcal{Z}_i$'s, we get the required contribution

$$|C_{e_R^*\tilde{\lambda}_i^+ H_u^0/H_d^0}| \equiv \frac{g_{YM} g_{T_B \bar{\mathcal{A}}_3} X^{T_B} \sim \mathcal{V}^{-\frac{16}{9}} \langle \mathcal{Z} \rangle \tilde{f}}{\sqrt{\hat{K}_{\mathcal{A}_3 \bar{\mathcal{A}}_3} \hat{K}^2_{\mathcal{Z}_1 \bar{\mathcal{Z}}_1}}} \mathcal{Z}_1 \bar{\chi}^{\bar{\mathcal{A}}_3} \tilde{\lambda}^0$$
$$\equiv \tilde{f}\left(\mathcal{V}^{-\frac{15}{9}} \frac{\langle \mathcal{Z}_i \rangle}{M_P}\right). \qquad (121)$$

Incorporating the results given in Eqs. (118)–(121) in Eq. (117), we have

$$|C_{e_L^* \chi_1^+ H_u^0/H_d^0}| \equiv \mathcal{V}^{-\frac{1}{10}}, \quad |C_{e_L^* \chi_2^+ H_u^0/H_d^0}| \equiv \mathcal{V}^{\frac{1}{10}},$$
$$|C_{e_R^* \chi_1^- H_u^0/H_d^0}| \equiv \tilde{f}\mathcal{V}^{-\frac{3}{2}}, \quad |C_{e_R^* \chi_2^- H_u^0/H_d^0}| \equiv \tilde{f}\mathcal{V}^{-\frac{15}{9}} \frac{\langle \mathcal{Z}_i \rangle}{M_P}. \qquad (122)$$

Now, the one-loop EDM of the electron in this case will be given as [64]

$$\left.\frac{d}{e}\right|_{\chi_i^\pm} = \sum_i \frac{m_{\chi_j^\pm} Q'_{e(i)}}{(4\pi)^2}$$
$$\times \left[\frac{1}{m^2_{H_i^0}} \mathrm{Im}\left((C_{e_L^* \chi_i^+ H_i^0} C^*_{e_R^* \chi_j^- H_i^0}) A\left(\frac{m^2_{\chi_i^\pm}}{m^2_{H_i^0}}\right)\right)\right], \qquad (123)$$

where $m_{\chi_j^\pm}$ and $m^2_{H_i^0}$ correspond to the mass eigenstates of the chargino and Higgs mass matrix. The effective charge for this loop diagram will be $Q'_{e(i)} = Q_e C_{\chi_i^+ \chi_i^- \gamma}$ where $C_{\chi_1^+ \chi_1^- \gamma} = C_{\tilde{H}_i^+ \tilde{H}_i^- \gamma}$, $C_{\chi_2^+ \chi_2^- \gamma} = C_{\tilde{\lambda}_i^+ \tilde{\lambda}_i^- \gamma}$. The contributions of both the Higgsino-Higgsino-gauge boson vertex and gaugino-gaugino-gauge boson have already been obtained in the context of $\mathcal{N} = 1$ gauged supergravity in [38]. Using the same,

$$C_{\chi_1^+ \chi_1^- \gamma} \equiv \tilde{f}\mathcal{V}^{-\frac{5}{18}}, \qquad C_{\chi_2^+ \chi_2^- \gamma} \equiv \tilde{f}\mathcal{V}^{-\frac{11}{18}}. \qquad (124)$$

Utilizing the results of the $C_{e^*_{L/R} \chi_i^\pm H_i^0}$ vertices given in (122) and the assumption that the value of the phase factor associated with these couplings are of $\mathcal{O}(1)$, $m_{\chi_1^\pm} = m_{H_2} = \mathcal{V}^{\frac{59}{72}} m_{\frac{3}{2}}$, $m_{\chi_2^\pm} = \mathcal{V}^{\frac{2}{3}} m_{\frac{3}{2}}$, and $m_{H_1} \sim 125$ GeV as given in Sec. II, $\sin\theta_h = (0, 1]$, $\phi_e = (0, \frac{\pi}{2}]$, and $A(\frac{m^2_{\chi_i^\pm}}{m^2_{H_i^0}}) \equiv \frac{m^2_{H_i^0}}{m^2_{\chi_i^\pm}}$ by using (33), we have

$$\left.\frac{d}{e}\right|_{\chi_i^\pm} \equiv \frac{1}{\sqrt{2}(4\pi)^2} (\mathcal{V}^{-\frac{1}{10}+\frac{1}{10}}) \times \frac{\tilde{f}\mathcal{V}^{-\frac{5}{18}}}{\mathcal{V}^{\frac{59}{72}} m_{\frac{3}{2}}}$$
$$\equiv \mathcal{O}(10^{-32})\ \text{cm}, \quad \text{for } \mathcal{V} = \mathcal{O}(1) \times 10^4. \qquad (125)$$

C. One-loop diagrams involving a gravitino and sGoldstino in the loop

1. Gravitino contribution

In this section, we estimate the EDM of the electron (quark) by considering the gravitino as a propagator in one-loop diagrams despite the fact that these are logarithmically divergent. The loop diagrams are given in Fig. 6. To get the numerical estimate of the EDM corresponding to these diagrams, we first need to determine the contribution of the relevant vertices in $\mathcal{N} = 1$ gauged supergravity. The same are evaluated as follows: In $\mathcal{N} = 1$ gauged supergravity, the gravitino-fermion-sfermion vertex will be given as $\mathcal{L}_{\tilde{G}-f-\tilde{f}} = -\frac{1}{2}\sqrt{2} e g_{ij} \partial_\mu \phi^i \chi^j \gamma^\mu \gamma^\nu \psi_\mu$. The physical ψ_μ-lepton(quark)-slepton(squark) vertex will be given as

$$|C_{\tilde{G} e_L \tilde{e}_L}| \equiv \frac{g_{\mathcal{A}_1 \bar{\mathcal{A}}_1}}{\sqrt{\mathcal{K}_{\mathcal{A}_1 \bar{\mathcal{A}}_1} \mathcal{K}_{\mathcal{A}_1 \bar{\mathcal{A}}_1}}} \partial_\mu \mathcal{A}_1 \chi^{\mathcal{A}_1} \gamma^\mu \gamma^\nu \psi_\mu$$
$$\equiv \partial_\mu \mathcal{A}_1 \chi^{\mathcal{A}_1} \gamma^\mu \gamma^\nu \psi_\mu,$$
$$|C_{\tilde{G} u_L \tilde{u}_L}| \equiv \frac{g_{\mathcal{A}_2 \bar{\mathcal{A}}_2}}{\sqrt{\mathcal{K}_{\mathcal{A}_2 \bar{\mathcal{A}}_2} \mathcal{K}_{\mathcal{A}_2 \bar{\mathcal{A}}_2}}} \partial_\mu \mathcal{A}_2 \chi^{\mathcal{A}_2} \gamma^\mu \gamma^\nu \psi_\mu$$
$$\equiv \partial_\mu \mathcal{A}_2 \chi^{\mathcal{A}_2} \gamma^\mu \gamma^\nu \psi_\mu. \qquad (126)$$

The contribution of the physical sfermion-sfermion-photon vertices have already been obtained in Sec. III A, and the values of the same are given as $|C_{\tilde{e}_L \tilde{e}_L \gamma}| \equiv \mathcal{V}^{\frac{44}{45}} \tilde{\mathcal{A}}_1 \partial_\mu \tilde{\mathcal{A}}_1 A^\mu$, $|C_{\tilde{u}_L \tilde{u}_L \psi_\mu}| \equiv \mathcal{V}^{\frac{53}{45}} \tilde{\mathcal{A}}_2 \partial_\mu \tilde{\mathcal{A}}_2 A^\mu$. The contribution of the fermion-sfermion-photino(γ) vertex in the context of $\mathcal{N} = 1$ gauged supergravity action will given by

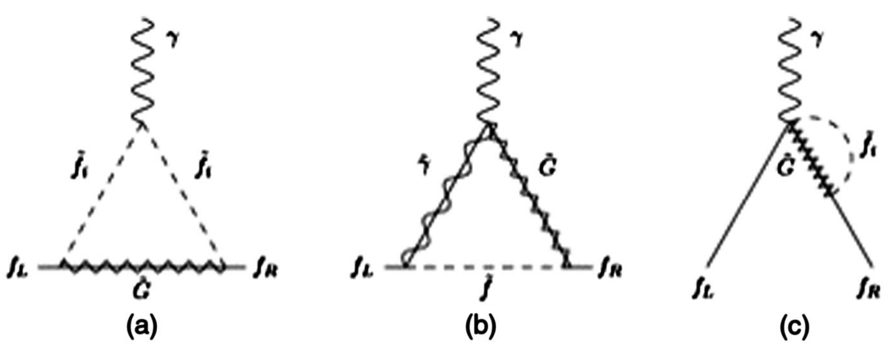

FIG. 6. One-loop diagrams involving a gravitino.

$\mathcal{L}_{f\tilde{f}\tilde{\gamma}} = g_{YM} g_{\mathcal{A}_I \bar{T}_B} X^{*B} \bar{\chi}^{\bar{\mathcal{A}}_I} \tilde{\gamma} + \partial_{\mathcal{A}_1} T_B D^B \bar{\chi}^{\bar{\mathcal{A}}_I} \tilde{\gamma}$. For $f = e$, by expanding $g_{a_1 \bar{T}_B}$ in the fluctuations linear in a_1 around its stabilized VEV, we have $g_{\bar{a}_1 T_B} = \mathcal{V}^{-\frac{2}{9}}(a_1 - \mathcal{V}^{-\frac{2}{9}})$ and $\partial_{\mathcal{A}_1} T_B \to \mathcal{V}^{\frac{10}{9}}(\mathcal{A}_1 - \mathcal{V}^{-\frac{2}{9}})$. Assuming that $g_{\bar{\mathcal{A}}_1 T_B} = \mathcal{O}(1) g_{\bar{a}_1 T_B}$ and using $X^{*B} = \kappa_4^2 \mu_7 Q_B$, $D^B = \frac{4\pi\alpha' \kappa_4^2 \mu_7 Q_B v^B}{\mathcal{V}}$ where $Q_B \sim \mathcal{V}^{\frac{1}{3}} \tilde{f}$ and $\kappa_4^2 \mu_7 \sim \frac{1}{\mathcal{V}}$, we get $|C_{e_L \tilde{e}_L \tilde{\gamma}}| \sim \frac{\mathcal{V}^{-\frac{2}{9}} \tilde{f}}{\sqrt{\mathcal{K}_{\mathcal{A}_1 \bar{\mathcal{A}}_1} \mathcal{K}_{\mathcal{A}_1 \bar{\mathcal{A}}_1}}} \tilde{\mathcal{A}}_1 \bar{\chi}^{\bar{\mathcal{A}}_1} \tilde{\gamma} \equiv \tilde{f} \mathcal{V}^{-1} \tilde{\mathcal{A}}_1 \bar{\chi}^{\bar{\mathcal{A}}_1} \tilde{\gamma}$. For $f = u$, using $g_{\bar{\mathcal{A}}_2 T_B} \sim \mathcal{O}(1) g_{\bar{a}_2 T_B} = \mathcal{V}^{-\frac{5}{4}}(a_2 - \mathcal{V}^{-\frac{1}{3}})$ and $\partial_{\mathcal{A}_2} T_B \to \mathcal{V}^{\frac{1}{9}}(\mathcal{A}_2 - \mathcal{V}^{-\frac{1}{3}})$, we have

$$|C_{u_L \tilde{u}_L \tilde{\gamma}}| \sim \frac{\mathcal{V}^{-\frac{11}{9}} \tilde{f}}{\sqrt{\mathcal{K}_{\mathcal{A}_2 \bar{\mathcal{A}}_2} \mathcal{K}_{\mathcal{A}_2 \bar{\mathcal{A}}_2}}} \tilde{\mathcal{A}}_2 \bar{\chi}^{\bar{\mathcal{A}}_2} \tilde{\gamma} \equiv \tilde{f} \mathcal{V}^{-\frac{4}{3}} \tilde{\mathcal{A}}_2 \bar{\chi}^{\bar{\mathcal{A}}_2} \tilde{\gamma}. \quad (127)$$

The contribution of the gravitino-fermion-sfermion-photon vertex in the context of $\mathcal{N} = 1$ gauged supergravity action will be given as $\mathcal{L} = -\frac{1}{2}\sqrt{2} e g_{\mathcal{A}_I T_B^*} X^B A_\mu \bar{\chi}^{\mathcal{A}_I} \gamma^\mu \gamma^\nu \psi_\mu$. Using the above-mentioned value of $g_{\mathcal{A}_1 T_B^*}$, $g_{\mathcal{A}_2 T_B^*}$, and X^B, the coefficient of the physical gravitino-lepton(quark)-slepton (squark)-photon vertex will be given as

$$\begin{aligned}
|C_{\tilde{G} e_L \tilde{e}_L \gamma}| &\sim \frac{\mathcal{V}^{-\frac{8}{9}} \tilde{f}}{\sqrt{\mathcal{K}_{\mathcal{A}_1 \bar{\mathcal{A}}_1} \mathcal{K}_{\mathcal{A}_1 \bar{\mathcal{A}}_1}}} A_\mu \bar{\chi}^{\mathcal{A}_1} \gamma^\mu \gamma^\nu \psi_\mu \\
&\equiv \tilde{f} \mathcal{V}^{-\frac{5}{3}} A_\mu \bar{\chi}^{\mathcal{A}_1} \gamma^\mu \gamma^\nu \psi_\mu, \\
|C_{\tilde{G} u_L \tilde{u}_L \gamma}| &\sim \frac{\mathcal{V}^{-\frac{35}{18}} \tilde{f}}{\sqrt{\mathcal{K}_{\mathcal{A}_2 \bar{\mathcal{A}}_2} \mathcal{K}_{\mathcal{A}_2 \bar{\mathcal{A}}_2}}} A_\mu \bar{\chi}^{\mathcal{A}_2} \gamma^\mu \gamma^\nu \psi_\mu \\
&\equiv \tilde{f} \mathcal{V}^{-\frac{5}{3}} A_\mu \bar{\chi}^{\mathcal{A}_2} \gamma^\mu \gamma^\nu \psi_\mu. \quad (128)
\end{aligned}$$

The contribution of the photon(γ)-photino($\tilde{\gamma}$)-gravitino(γ) vertex will be given $\mathcal{L} = \frac{i}{4} e \bar{\gamma}^\mu \lambda [\partial, \mathcal{A}] \psi_\mu$. We notice that there is no moduli space-dependent factor coming from this vertex.

The above Feynman diagrams involving a gravitino in a loop have been explicitly worked out in [65] to calculate the magnetic moment of the muon in the context of spontaneously broken minimal $\mathcal{N} = 1$ gauged supergravity. We explicitly utilize their results in a modified form to get the estimate of the EDM of the electron/quark in the $\mathcal{N} = 1$ gauged supergravity. The modified results of the magnetic moment of the electron after multiplying with volume-suppression factors coming from relevant vertices as calculated in Eqs. (126)–(128) are as follows.

For Fig. 6(a):

$$\begin{aligned}
a_f^{\mathrm{div}}|_{6(a)} \equiv{}& \tilde{f} \mathcal{V}^a (G_N m_f^2/\pi) \sum_{j=1,2} \left[\Gamma(\epsilon - 1) \left[-\frac{1}{90}\mu^2 + \frac{1}{18}\mu_j^2 \right] \right. \\
&\left. + \Gamma(\epsilon) \left[\frac{2}{45}\mu^2 + \frac{2}{9} \right] + (-1)^j \sin\theta \Gamma(\epsilon - 1)[-\mu_j^2/3\mu] \right], \quad (129)
\end{aligned}$$

where $\mathcal{V}^a$ is the Calabi-Yau volume-suppression factor. Here, $\mu = m_f/m_{\frac{3}{2}}$ and $\mu_j = m_{\tilde{f}_j}/m_{3/2}$, $j = 1, 2$, m is the lepton mass, and $m_{\frac{3}{2}}$ is the gravitino mass. $m_{\tilde{f}_1}$ and $m_{\tilde{f}_2}$ are the eigenvalues of the diagonalized slepton (squark) mass matrix. In our setup, $\sin\theta = 1$. Using $m_{\tilde{f}_1} = m_{\tilde{f}_2} = \mathcal{V}^{\frac{1}{2}} m_{\frac{3}{2}}$, $m_{\frac{3}{2}} = \mathcal{V}^{-2} M_P$, and $m_e = \mathcal{O}(1)$ MeV, we have $\mu_1 = \mu_2 = \frac{1}{\mathcal{V}}$ and $\mu = 10^{-11}$ for $\mathcal{V} = 10^5$.

For $f = e$, incorporating these values, the dominant contribution will be of the form

$$\begin{aligned}
a_e^{\mathrm{div}}|_{6(a)} &\equiv \tilde{f} \mathcal{V}^{\frac{44}{45}} (G_N m_e^2/\pi) \left[\frac{1}{18\mathcal{V}^2} \Gamma(\epsilon - 1) + \frac{2}{9} \Gamma(\epsilon) \right] \\
&\equiv \tilde{f} \mathcal{V}^{\frac{44}{45}} (G_N m_e^2/\pi) \left[\frac{1}{18\mathcal{V}^2} \Gamma(\epsilon - 1) + \frac{1}{18\mathcal{V}^2} \Gamma(\epsilon) + a' \right], \quad (130)
\end{aligned}$$

where $a' = (\frac{2}{9} - \frac{1}{18\mathcal{V}^2})\Gamma(\epsilon)$ is the divergent piece. Using $-\Gamma(\epsilon - 1) = \Gamma(\epsilon)(1 + \epsilon)$, the finite contribution will be given as $a_e^{\mathrm{finite}}|_{6(a)} \equiv \frac{1}{18} \tilde{f} \mathcal{V}^{-\frac{46}{45}} (G_N m_e^2/\pi)$. Similarly, using the volume-suppression factor coming from quark-quark-photon vertex, we get $a_u^{\mathrm{finite}}|_{6(a)} \equiv \frac{1}{18} \tilde{f} \mathcal{V}^{-\frac{37}{45}} (G_N m_u^2/\pi)$. Now we use the relation between the anomalous magnetic moment and the electric dipole moment to get the numerical estimate of the EDM of the electron in this case. As given in [66], $a_f = \frac{2|m_f|}{eQ_f} |d_f| \cos\phi$, where m_f and Q_f correspond to the mass and charge of the fermion; d_f is the electric dipole moment of the fermion, and ϕ is defined as $\phi \equiv \arg(d_f m_f^*)$. We consider that in the loop diagrams involving sfermions as propagators, the nontrivial phase responsible to generate the EDM appears from eigenstates of the sfermion mass matrix (off-diagonal component of the slepton mass matrix), and we assume the value of the same as $\phi_{d_f} \exists (0, \frac{\pi}{2}]$. The first-generation electron/quark mass has been calculated from the complex effective Yukawa coupling ($\mathcal{Y}^{\mathrm{eff}}_{\mathcal{Z}_I \mathcal{A}_{1/3} \mathcal{A}_{2/4}}$) in $\mathcal{N} = 1$ gauged supergravity, and there is a distinct phase factor ϕ_{y_e/y_u} associated with the same. Using the fact that $\phi_{d_f} \neq \phi_{y_e/y_u}$, the relative phase between the two will be in the interval $\phi \exists (0, \frac{\pi}{2}) \sim \mathcal{O}(1)$. Hence,

$$\begin{aligned}
\left.\frac{d_e}{e}\right|_{6(a)} &= 2|m_e| a_e^{\mathrm{finite}}|_{6(a)} \equiv \frac{1}{18} \tilde{f} \mathcal{V}^{-\frac{46}{45}} (G_N m_e/\pi) \equiv 10^{-67}\ \mathrm{cm}, \\
\left.\frac{d_u}{e}\right|_{6(a)} &= 2|m_u| a_u^{\mathrm{finite}}|_{6(a)} \equiv \frac{1}{18} \tilde{f} \mathcal{V}^{-\frac{37}{45}} (G_N m_u/\pi) \equiv 10^{-67}\ \mathrm{cm}. \quad (131)
\end{aligned}$$

For Fig. 6(b):

$$a_f^{\rm div}|_{6(b)} \equiv \tilde{f}\mathcal{V}^{-a}(G_N m_f^2/\pi)\sum_{j=1,2}\Bigg[\Gamma(\epsilon-1)\left[\frac{1}{20}\mu^2-\frac{1}{6}\mu_j^2\right] + \Gamma(\epsilon)\left[-\frac{7}{60}\mu^2\right]+(-1)^j\sin 2\alpha\Gamma(\epsilon-1)[\mu_j^2/\mu]\Bigg],$$

where $f = e, u$.

For $f = e$, incorporating the values of the masses and simplifying, now we will have

$$\begin{aligned} a_e^{\rm div}|_{6(b)} &\equiv \tilde{f}\mathcal{V}^{-1}(G_N m_e^2/\pi)\left[-\frac{1}{6\mathcal{V}^2}\Gamma(\epsilon-1)-\frac{7}{60}\mu^2\Gamma(\epsilon)\right] \\ &\equiv \tilde{f}\mathcal{V}^{-1}(G_N m_f^2/\pi)\left[-\frac{7}{60}\mu^2\Gamma(\epsilon-1)-\frac{7}{60}\mu^2\Gamma(\epsilon)+a'\right], \end{aligned} \tag{132}$$

where $a' = (-\frac{1}{6\mathcal{V}^2}+\frac{7}{60}\mu^2)\Gamma(\epsilon-1)$ is the divergent piece. Picking up the finite contribution, we get $a_e^{\rm finite}|_{6(b)} \equiv \frac{7}{60}\tilde{f}\mathcal{V}^{-1}(G_N m_f^2/\pi)\mu^2$, and, therefore,

$$\frac{d_e}{e}\bigg|_{6(b)} \equiv 2|m_e|\, a_e^{\rm finite}|_{6(b)} \equiv 10^{-65}\ {\rm GeV}^{-1} \equiv 10^{-79}\ {\rm cm}. \tag{133}$$

Similarly, using the volume-suppression factor coming from the quark-quark-photon vertex,

$$\frac{d_u}{e}\bigg|_{6(b)} \equiv 10^{-64}\ {\rm GeV}^{-1} \equiv 10^{-78}\ {\rm cm}. \tag{134}$$

For Fig. 6(c):

$$a_f^{\rm div}|_{6(c)} \equiv \tilde{f}\mathcal{V}^{-\frac{5}{3}}(G_N m_f^2/\pi)\sum_{j=1,2}\Bigg[\Gamma(\epsilon-1)\left[-\frac{1}{90}\mu^2+\frac{1}{9}\mu_j^2\right] + \Gamma(\epsilon)\left[\frac{1}{10}\mu^2-\frac{2}{9}\right] + (-1)^j\sin 2\alpha\Gamma(\epsilon-1)[-2\mu_j^2/3\mu]\Bigg].$$

As similar to the above, incorporating the value of the masses and further simplifying, the dominant contribution is given by

$$\begin{aligned} a_f^{\rm div}|_{6(c)} &\equiv \tilde{f}\mathcal{V}^{-\frac{5}{3}}(G_N m_f^2/\pi)\left[\frac{1}{9\mathcal{V}^2}\Gamma(\epsilon-1)-\frac{2}{9}\Gamma(\epsilon)\right] \\ &\equiv \tilde{f}\mathcal{V}^{-\frac{5}{3}}(G_N m_f^2/\pi)\left[\frac{1}{9\mathcal{V}^2}\Gamma(\epsilon-1)+\frac{1}{9\mathcal{V}^2}\Gamma(\epsilon)+a'\right], \end{aligned} \tag{135}$$

where $a' = (-\frac{2}{9}-\frac{1}{9\mathcal{V}^2})\Gamma(\epsilon)$ is the divergent piece. Considering the finite piece, $a_q^{\rm finite}|_{6(c)} = \frac{1}{9}\tilde{f}\mathcal{V}^{-\frac{11}{3}}(G_N m_f^2/\pi)$. Again using $\frac{d_f}{e}|_{6(c)} = 2|m_f|a_e^{\rm finite}|_{6(c)}$, we get

$$\frac{d_e}{e}\bigg|_{6(c)} = \frac{d_u}{e}\bigg|_{6(c)} \equiv 10^{-66}\ {\rm GeV}^{-1} \equiv 10^{-80}\ {\rm cm}. \tag{136}$$

Hence, the overall contribution of the EDM of the electron as well as the neutron/quark in the case of one-loop Feynman diagrams involving a gravitino is

$$\frac{d_e}{e}\bigg|_{\tilde{G}} = \frac{d_n}{e}\bigg|_{\tilde{G}} = \frac{d_{e/u}}{e}\bigg|_{6(a)} + \frac{d_{e/u}}{e}\bigg|_{6(b)} + \frac{d_{e/u}}{e}\bigg|_{6(c)} \equiv 10^{-67}\ {\rm cm}. \tag{137}$$

2. sGoldstino contribution

In supersymmetric models, the sGoldstino is the bosonic component of the superfield corresponding to which there is an F-term (D-term) supersymmetry breaking. In our setup, supersymmetry is broken in the bulk sector, and the scale of the same is governed by the F term (assuming that in the dilute flux approximation $V_D \ll V_F$) corresponding to bulk fields $(F^{\tau_S}, F^{\tau_B}, \mathcal{G}^a)$ where τ_S and τ_B correspond to small and big divisor volume moduli, and $\mathcal{G}^a$ corresponds to complexified NS-NS and RR axions. It was shown in [38], at M_s, $|F^{\tau_S}| > |F^{\mathcal{G}^a}|$, $|F^{\tau_B}|$. From Sec. IVA, the requirement of the quark-quark-photon coupling to be the SM at the EW scale, we see that $|F^{\tau_B}|$ is the most dominant F term at the EW scale. To obtain an estimate of the off-shell Goldstino multiplet, we consider the same to be (τ_B, χ_B, F^B), where τ_B is a complex scalar field. Here, we identify σ_B with the scalar (sGoldstino) field and ρ_B with the pseudoscalar (sGoldstino) field.

3. Mass of the sGoldstino

The dominant contribution to the F-term potential at the string scale M_s is given by $V = \|F^{\tau_S}\|^2$, where $F^{\bar{\tau}_S} = e^{K/2}\bar{\partial}^{\bar{\tau}_S}\partial^{\beta}KD_{\beta}W$.[5] At the EW scale, the F-term potential receives the dominant contribution from the $\|D_{\tau_B}W\|^2$ term and is estimated to be $V(n_s=2)|_{\rm EW} \sim e^K K^{\tau_S\bar{\tau}_B}D_{\tau_S}WD_{\bar{\tau}_B}\bar{W} + e^K K^{\tau_B\bar{\tau}_B}|D_{\tau_B}W|^2$, near $\langle\sigma_S\rangle \sim \frac{\ln\mathcal{V}}{(\mathcal{O}(1))^4_{\sigma_S}}$, $\langle\sigma_B\rangle \equiv \frac{\mathcal{V}^{\frac{2}{3}}}{(\mathcal{O}(1)_{\sigma_B})^4}$ yields

$$\frac{\partial^2 V}{\partial\sigma_B^2}\bigg|_{\rm EW} \equiv \mathcal{V}^{-\frac{1}{3}}m_{3/2}^2(\mathcal{O}(1)_{\sigma_B})^2\left((\mathcal{O}(1)_{\sigma_S})^2 + \frac{(\mathcal{O}(1)_{\sigma_S})^6}{\ln\mathcal{V}}\right).$$

For the aforementioned $\mathcal{O}(1)_{\sigma_B} = \frac{\mathcal{O}(1)_{\sigma_S}}{2} \equiv 3.5$ for $\mathcal{V} \sim 10^4$, $(\frac{\partial^2 V}{\partial\sigma_B^2})_{\rm EW} \equiv \mathcal{V}^{\frac{4}{3}}$, and the canonically normalized coefficient

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[5]We note that $e^{K(\tau_{S,B},G^a,z^i,a_I,\ldots)}\bar{\partial}^{\bar{\mathcal{I}}}\partial^{\mathcal{J}}K(\tau_{S,B},G^a,z^i,a_I,\ldots)\times D_{\mathcal{J}}WD_{\bar{\mathcal{I}}}\bar{W}(\mathcal{I}\equiv T_{s,b},G^a,z^i,a_I,\ldots) = e^{K(\tau_{z,b},G^a,z^i,a_I,\ldots)}\bar{\partial}^{\bar{\alpha}}\partial^{\beta}K(\tau_{S,B}, G^a,z^i,a_I,\ldots)D_{\beta}WD_{\bar{\alpha}}\bar{W}(\beta=\tau_{S,B},G^a,z^i,a_I)$; however, $G_{\mathcal{I}\bar{\mathcal{J}}} = \partial_{\mathcal{I}}\bar{\partial}_{\bar{\mathcal{J}}}K(T_{S,B},G^a,z^i,a_I,\ldots)$, $G_{\alpha\bar{\beta}} \neq \partial_{\alpha}\bar{\partial}_{\bar{\beta}}K(\tau_{S,B},G^a,z^i,a_I,\ldots)$ as $\tau_{S,B}$ is not an $\mathcal{N}=1$ chiral coordinate.

quadratic in the fluctuations yields the sGoldstino mass estimate:

$$m_{\tau_B} \sim \sqrt{\frac{\left(\frac{\partial^2 V}{\partial \sigma_B^2}\right)_{\rm EW}}{\kappa_4^2 \mu_7 K^{\tau_B \bar{\tau}_B}_{\rm EW}}} \sim \mathcal{O}(1) m_{3/2}.$$

It will be interesting to get the contribution of the same to the electron/neutron EDM.

To get the analysis of one-loop diagrams involving the sGoldstino, we consider only the scalar sGoldstino field (σ_s) and first calculate the contribution of the vertices involving the sGoldstino in the context of $\mathcal{N}=1$ gauged supergravity. The coefficient of the lepton(e_L)-scalar[sGoldstino(σ_B)]-lepton(e_R) vertex has been calculated by expanding $\frac{e^{\frac{K}{2}}}{2}(\mathcal{D}_{\mathcal{A}_1} D_{\mathcal{A}_3} W)\bar{\chi}^{\mathcal{A}_1}\chi^{\mathcal{A}_3}$ in the fluctuations linear in σ_B in $\mathcal{N}=1$ gauged supergravity. By expanding the above in the fluctuations linear in $\sigma_B \to \sigma_B + \mathcal{V}^{\frac{2}{3}} M_P$, on simplifying, we have $\frac{e^{\frac{K}{2}}}{2}\mathcal{D}_{\mathcal{A}_1} D_{\mathcal{A}_3} W \chi^{\mathcal{A}_1}\bar{\chi}^{\mathcal{A}_3} \sim \mathcal{V}^{-\frac{13}{3}} \delta\sigma_B \chi^{\mathcal{A}_1}\bar{\chi}^{\mathcal{A}_3}$. The physical lepton($e_L$)-sGoldstino($\sigma_B$)-lepton($e_R$) vertex will be given as

$$|C_{\delta\sigma_B e_L e_R^c}| \equiv \frac{\mathcal{V}^{-\frac{13}{3}}}{\sqrt{k_4^2 \mu_7 G^{\tau_B \bar{\tau}_B} \hat{K}_{\mathcal{A}_1 \bar{\mathcal{A}}_1} \hat{K}_{\mathcal{A}_3 \bar{\mathcal{A}}_3}}} \equiv \mathcal{V}^{-\frac{92}{15}}, \quad \text{for } \mathcal{V} \sim 10^5. \tag{138}$$

Similarly, the coefficient of the quark(u_L)-scalar[sGoldstino(σ_B)]-lepton(u_R) vertex can be calculated by expanding $\frac{e^{\frac{K}{2}}}{2}(\mathcal{D}_{\mathcal{A}_2} D_{\mathcal{A}_4} W)\bar{\chi}^{\mathcal{A}_2}\chi^{\mathcal{A}_4}$ in the fluctuations linear in σ_B in $N=1$ gauged supergravity. Using the similar procedure, we get $\frac{e^{\frac{K}{2}}}{2}\mathcal{D}_{\mathcal{A}_2} D_{\mathcal{A}_4} W \chi^{\mathcal{A}_2}\bar{\chi}^{\mathcal{A}_4} \sim \mathcal{V}^{-4} \delta\sigma_B \chi^{\mathcal{A}_2}\bar{\chi}^{\mathcal{A}_4}$. Therefore, the physical quark(u_L)-sGoldstino(σ_B)-quark(u_R) vertex will be given as

$$|C_{\delta\sigma_B u_L u_R^c}| \sim \frac{\mathcal{V}^{-5}}{\sqrt{k_4^2 \mu_7 G^{\tau_B \bar{\tau}_B} \hat{K}_{\mathcal{A}_2 \bar{\mathcal{A}}_2} \hat{K}_{\mathcal{A}_4 \bar{\mathcal{A}}_4}}} \equiv \mathcal{V}^{-\frac{33}{5}}, \quad \text{for } \mathcal{V} \sim 10^5. \tag{139}$$

In $\mathcal{N}=1$ supergravity, the contribution of the photon-sGoldstino(scalar)-photon will be accommodated by the gauge kinetic term $\mathcal{L} = \mathrm{Re}(T_B) F \wedge *_4 F$, where $\mathrm{Re}(T_B) = \sigma_B - C_{i\bar{j}} a_i a_{\bar{j}}$. Considering $\sigma_B \to \langle\sigma_B\rangle + \delta\sigma_B$, the coefficient of the physical vertex will be given as $|C_{\gamma\gamma\delta\sigma_B}| \equiv \frac{1/M_p}{\sqrt{k_4^2 \mu_7 G^{\tau_B \bar{\tau}_B}}} \sim \frac{\mathcal{V}^{-\frac{2}{3}}}{M_p}$. The possibility of getting the fermion-fermion-photon vertex $C_{ff^*\gamma} \equiv \mathcal{O}(1)$ has been shown in Sec. III A.

Now we use the values of the coefficients of the relevant vertices to evaluate the estimate of the EDM for the loop diagrams given in Figs. 7(a) and 7(b). The diagrams have

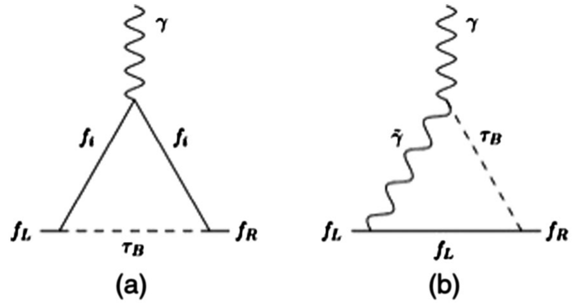

FIG. 7. One-loop diagrams involving an sGoldstino.

been evaluated in [67] to determine the estimate of the magnetic moment of the muon in $\mathcal{N}=1$ global supersymmetry. Utilizing their results in a modified form in the context of $\mathcal{N}=1$ gauged SUGRA and the relation between the magnetic moment and EDM as given above, for Fig. 7(a), the magnitude of the electric dipole moment will be

$$\left|\frac{d_f}{e}\right|_{7(a)} = \frac{m_f}{16\pi^2}\cos\phi\left[(C_{\delta\sigma_B f_L f_R^c})^2 \int_0^1 dx \frac{x^2(2-x)}{m_{\sigma_B}^2(1-x) + m_f^2 x^2}\right]. \tag{140}$$

Putting the value of $|C_{\delta\sigma_B e_L e_R^c}| \equiv \mathcal{V}^{-\frac{92}{15}}$, $|C_{\delta\sigma_B u_L u_R^c}| \equiv \mathcal{V}^{-\frac{33}{6}}$, and the value of masses $m_{\sigma_B} = m_{\frac{3}{2}}$, $m_e = 0.5$ MeV, we get

$$\left|\frac{d_e}{e}\right|_{7(a)} \equiv 10^{-95}\ \text{cm} \quad \text{and} \quad \left|\frac{d_n}{e}\right|_{7(a)} \equiv 10^{-89}\ \text{cm}. \tag{141}$$

For Fig. 7(b):

$$\left|\frac{d_f}{e}\right|_{7(b)} = \frac{C_{\delta\sigma_B f_L f_R^c} C_{\gamma\gamma\delta\sigma_B}}{8\pi^2}\left[\Delta_{UV} - \frac{1}{2} - \int_0^1 dx \int_0^{1-x} dy \log\left[\frac{m_{\sigma_B}^2 y + m_f^2 x^2}{\mu^2}\right]\right], \tag{142}$$

where $\Delta_{UV} = \log[\frac{\Delta_{\rm UV}^2}{\mu^2}] - 1$. Incorporating the values of the relevant inputs and considering the finite piece,

$$\left|\frac{d_e}{e}\right|_{7(b)} \equiv 10^{-72}\ \text{cm}, \qquad \left|\frac{d_n}{e}\right|_{7(b)} \equiv 10^{-68}\ \text{cm}. \tag{143}$$

Hence, the overall contribution of the sGoldstino to the EDM of the electron/neutron is

$$\left|\frac{d_e}{e}\right|_{\rm sGoldstino} = \left|\frac{d_e}{e}\right|_{7(a)} + \left|\frac{d_e}{e}\right|_{7(b)} \equiv 10^{-72}\ \text{cm}, \tag{144}$$

$$\left|\frac{d_n}{e}\right|_{\rm sGoldstino} = \left|\frac{d_n}{e}\right|_{7(a)} + \left|\frac{d_n}{e}\right|_{7(b)} \equiv 10^{-68}\ \text{cm}. \tag{145}$$

TABLE II. Results of the EDM of the electron/neutron for all possible one-loop diagrams.

One-loop particle exchange	Origin of complex phase	$d_e(e\ \mathrm{cm})$	$d_n(e\ \mathrm{cm})$
$\lambda^0\tilde{f}$	Diagonalized sfermion mass eigenstates	10^{-39}	10^{-38}
$\chi_i^0\tilde{f}$	”	10^{-37}	10^{-34}
$f\tilde{f}$	”	10^{-45}	10^{-45}
fh_i^0	Digonalized Higgs mass eigenstates	10^{-34}	10^{-33}
$\chi^\pm h_i^0$	”	10^{-32}	…
Gravitino$(\tilde{G})\,\tilde{f}$	Diagonalized sfermion mass eigenstates	10^{-67}	10^{-67}
sGoldstino $\tilde{f}$	”	10^{-72}	10^{-68}

The results of all possible one-loop diagrams contributing to the EDM of the electron/neutron are summarized in Table II.

V. TWO-LOOP-LEVEL BARR-ZEE-TYPE CONTRIBUTION TO THE ELECTRIC DIPOLE MOMENT

In the two-loop diagrams discussed in this section, the CP-violating effects are mainly demonstrated by complex effective Yukawa couplings which include R-parity violating couplings, SM-like Yukawa couplings, as well as couplings involving a Higgsino, and complex scalar trilinear couplings in the context of $\mathcal{N}=1$ gauged supergravity. In the subsection given below, we present the contribution of individual Barr-Zee-type diagrams formed by including an internal fermion loop generated by R-parity violating interactions, SM-like Yukawa interactions, and gaugino(gaugino)-Higgsino(Higgsino)-Higgs couplings. The two-loop diagrams are shown in Fig. 8.

A. Two-loop-level Barr-Zee Feynman diagrams involving an internal Fermion loop

1. Higgs contribution

For the two-loop diagram given in Fig. 8(a), the interaction Lagrangian is governed by Yukawa couplings given as

$$\mathcal{L} \supset \hat{Y}_{H_i^0 u_L u_R^c} H_i^0 u_{kL} u_{kR}^c + \hat{Y}^*_{H_i^0 e_L e_R^c} H_i^0 e_{jL} e_{jR}^c + \mathrm{H.c.} \qquad (146)$$

We have already given the estimate of effective Yukawa couplings for the first generation of leptons and quarks in [38] in the context of $\mathcal{N}=1$ gauged supergravity. Using those results, we have

$$\hat{Y}_{H_i^0 e_L e_R^c} \sim \hat{Y}^{\mathrm{eff}}_{\mathcal{Z}_i \mathcal{A}_1 \mathcal{A}_3} \equiv \mathcal{V}^{-\frac{47}{45}} e^{i\phi_{y_e}},$$
$$\hat{Y}_{H_i^0 u_L u_R^c} \sim \hat{Y}^{\mathrm{eff}}_{\mathcal{Z}_i \mathcal{A}_2 \mathcal{A}_4} \equiv \mathcal{V}^{-\frac{17}{18}} e^{i\phi_{y_u}} \quad \text{for } \mathcal{V} = 105, \qquad (147)$$

where $e^{i\phi_{y_e}}$ and $e^{i\phi_{y_u}}$ are nonzero phases of the aforementioned Yukawa couplings.

For a two-loop Barr-Zee diagram involving an internal fermion loop and taking into account the chirality flip between the internal loop and external line, the analytical expression has been derived in [68,69]. Using the same, the

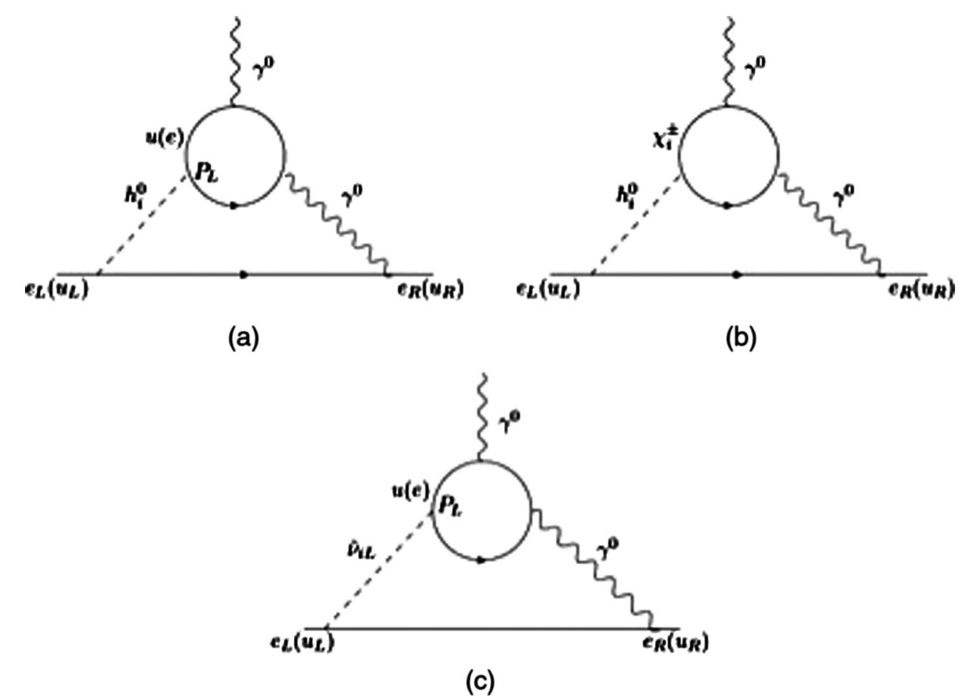

FIG. 8. Two-loop diagrams involving fermions in the internal loop.

electric dipole moment of the electron for the loop diagram given in Fig. 8(a) will be[6]

$$\frac{d}{e}\bigg|_H = \sum_{i=1,2} \mathrm{Im}(\hat{Y}_{H_i^0 e_L e_R^c} \hat{Y}_{H_i^0 u_L u_R^c}) \frac{\alpha_{em} Q_u^2 Q_e}{16\pi^3 m_{ej}} (f(z_1) - g(z_1)), \quad (148)$$

and the EDM of the neutron will be given as

$$\frac{d}{n}\bigg|_H = \sum_{i=1,2} \mathrm{Im}(\hat{Y}_{H_i^0 e_L e_R^c} \hat{Y}_{H_i^0 u_L u_R^c}) \frac{\alpha_{em} Q_e^2 Q_u}{16\pi^3 m_{ej}} (f(z_2) - g(z_2)), \quad (149)$$

where $z_1 = \frac{m_e^2}{m^2_{H_i^0}}$, $z_2 = \frac{m_u^2}{m^2_{H_i^0}}$, and

$$f(z) = \frac{z}{2}\int_0^1 dx \frac{1-2x(1-x)}{x(1-x)-z} \ln\left(\frac{x(1-x)}{z}\right),$$
$$g(z) = \frac{z}{2}\int_0^1 dx \frac{1}{x(1-x)-z} \ln\left(\frac{x(1-x)}{z}\right). \quad (150)$$

Using the value of masses $m_{H_1^0} = 125$ GeV, $m_{H_2^0} = \mathcal{V}^{\frac{59}{72}} m_{\frac{3}{2}}$, and $m_e = 0.5$ GeV, $f(\frac{m_e^2/m_u^2}{m^2_{H_1^0}}) = g(\frac{m_e^2/m_u^2}{m^2_{H_1^0}}) = 10^{-10}$, $f(\frac{m_e^2/m_u^2}{m^2_{H_2^0}}) = g(\frac{m_e^2/m_u^2}{m^2_{H_2^0}}) = 10^{-23}$. Utilizing the same and assuming $e^{i(\phi_{y_e} - \phi_{y_u})} = (0,1]$, Eqs. (148) and (149) reduce to give the EDM result as follows:

$$\frac{d}{e}\bigg|_H = \frac{d}{n}\bigg|_H \sim \mathcal{V}^{-2} \times 10^{-2} \times 10^{-10}$$
$$= 10^{-22}\ \mathrm{GeV}^{-1} \equiv 10^{-36}\ \mathrm{cm}. \quad (151)$$

2. Chargino contribution

In the loop diagram [Fig. 8(b)], the general Lagrangian governing the interaction of the charginos will be

$$\mathcal{L} \supset C_{ikk} H_i^0 \chi_{kL}^+ \chi_{kR}^- + \hat{Y}^*_{H_i^0 e_L e_R^c} H_i^0 e_{jL} e_{jR}^c + \hat{Y}^*_{H_i^0 u_L u_R^c} H_i^0 u_{jL} u_{jR}^c + \mathrm{H.c.} \quad (152)$$

We evaluate the contribution of the chargino($\chi_i^\pm$)-Higgs-chargino($\chi_1^\pm$) vertex in $\mathcal{N}=1$ gauged supergravity. As described in the Appendix, $\chi_1^\pm$ and $\chi_2^\pm$ correspond to a Higgsino ($\tilde{H}_i^\pm$) with a very small admixture of gaugino ($\lambda_i^\pm$) and vice versa. So, $C_{\chi_i^+\chi_1^- H_i^0} \equiv C_{\tilde{H}_i^+ \tilde{H}_i^- H_i^0}$ and $C_{\chi_2^+\chi_2^- H_i^0} \equiv C_{\tilde{\lambda}_i^+ \tilde{\lambda}_i^- H_i^0}$.

[6]We consider $Q_e' = C_{ee^*\gamma} Q_e \sim Q_e$ because $C_{ee^*\gamma} \sim \mathcal{O}(1)$ as shown in Sec. IV C. Similarly, $Q_u' \sim Q_u$.

Higgsino(χ_{kL}^-)-Higgs-Higgsino(χ_{kR}^+) vertex.—Given that the Higgsino is a majorana particle, $\chi_{kR}^+ = (\chi_{kL}^-)^c$. In our model, the Higgsino has been identified with position moduli $\mathcal{Z}_i$; the contribution of this vertex in $\mathcal{N}=1$ gauged sypergravity will be given by expanding $e^{\frac{K}{2}} \mathcal{D}_{\mathcal{Z}_i} \mathcal{D}_{\bar{\mathcal{Z}}_i} W$ in the fluctuations linear in $\mathcal{Z}_i$ about its stabilized VEV. Since $SU(2)_L$ symmetry is not conserved for this vertex, we will expand the above in the fluctuations quadratic in $\mathcal{Z}_i$, giving a VEV to one of the $\mathcal{Z}_i$. Considering $z_i \to \mathcal{V}^{\frac{1}{18}} + \delta z_i$, we have $\mathcal{D}_{z_i} \mathcal{D}_{\bar{z}_i} W = \mathcal{V}^{-\frac{16}{9}} z_i \langle z_i \rangle$. Using $\mathcal{D}_{\mathcal{Z}_i} \mathcal{D}_{\bar{\mathcal{Z}}_i} W \sim \mathcal{D}_{\bar{z}_1} D_{z_i} W$, the physical vertex will be given as $C_{\chi_i^+ \chi_1^- H_i^0} \equiv C_{\tilde{H}_i^+ \tilde{H}_i^- H_i^0} = \frac{\mathcal{V}^{-\frac{7}{4}}}{(\sqrt{\hat{K}_{\mathcal{Z}_1 \bar{\mathcal{Z}}_1}})^4} = \mathcal{V}^{\frac{1}{4}} e^{i\phi_{\chi_1}}$, where ϕ_{χ_1} corresponds to the nonzero phase associated with the aforementioned coupling.

Gaugino(λ_{kR}^+)-Higgs-gaugino(λ_{kL}^+) vertex.—The coefficient of this vertex will be given from the kinetic term of the gaugino. The interaction term corresponding to this coupling will be given by considering term $\mathcal{L} = i\bar{\lambda}_L \gamma^m \frac{1}{4} (K_{\mathcal{Z}_i} \partial_m \mathcal{Z}_i - \mathrm{c.c.}) \lambda_L + (\partial_{\mathcal{Z}_i} T_B) \bar{\lambda}_L \gamma^m \frac{1}{4} (K_{\mathcal{Z}_i} \partial_m \mathcal{Z}_i - \mathrm{c.c.}) \lambda_L$, where λ_L corresponds to the gaugino. Given that charged (gauginos) are either $SU(2)_L$ singlets or triplets, the aforementioned vertex does not preserve $SU(2)_L$ symmetry—one has to obtain the term bilinear in $\mathcal{Z}_i$ such that we give a VEV to one of the $\mathcal{Z}_i$'s. Since $(\partial_{\mathcal{Z}_i} T_B)$ does not contain terms bilinear in $\mathcal{Z}_i$, which are needed to ensure $SU(2)_L$ symmetry, the second term contributes zero to the given vertex. In terms of the undiagonalized basis, $\partial_{z_i} K \sim \mathcal{V}^{-\frac{2}{3}} \langle z_i \rangle$, and using $\partial_{\mathcal{Z}_i} K \sim \mathcal{O}(1) \partial_{z_i} K$, we have $\partial_{\mathcal{Z}_i} K \sim \mathcal{V}^{-\frac{2}{3}} \langle \mathcal{Z}_i \rangle$. Incorporating the same, we get

$$\mathcal{L} = \frac{\mathcal{V}^{-\frac{2}{3}} \langle \mathcal{Z}_i \rangle \bar{\lambda}_L \frac{\not{\partial}\mathcal{Z}_i}{M_P} \lambda_L}{\sqrt{(\hat{K}_{\mathcal{Z}_1 \bar{\mathcal{Z}}_1})^2}} \sim \mathcal{V}^{\frac{13}{36}} h \bar{\lambda}_L \frac{\not{p}_h}{M_P} \lambda_L$$
$$\sim \mathcal{V}^{\frac{13}{36}} h \bar{\lambda}_L \frac{\gamma \cdot (p_{e_L} + p_{e_R})}{M_P} \lambda_L. \quad (153)$$

Therefore, $C_{\chi_2^+ \chi_2^- H_i^0} \equiv C_{H_i^0 \lambda_R^- \lambda_L^+} \sim \mathcal{V}^{\frac{13}{36}} \frac{m_e}{M_P} e^{i\phi_{\tilde{\lambda}_1^0}}$ where $\phi_{\tilde{\lambda}_1^0}$ corresponds to the nonzero phase associated with the aforementioned coupling.

The contribution of the gaugino-gaugino-gauge boson as well as the Higgsino-Higgsino-gauge boson have been already evaluated in the context of $\mathcal{N}=1$ gauged supergravity. The volume-suppression factors corresponding to these vertices are as follows:

$$|C_{\chi_1^+ \chi_1^- \gamma}| \equiv |C_{\tilde{H}_u^+ \tilde{H}_d^- \gamma}| \equiv \tilde{f} \mathcal{V}^{-\frac{5}{18}};$$
$$|C_{\chi_2^+ \chi_2^- \gamma}| \equiv |C_{\tilde{\lambda}_i^+ \tilde{\lambda}_i^- \gamma}| \equiv \tilde{f} \mathcal{V}^{-\frac{11}{18}}. \quad (154)$$

Now, the EDM of the electron for the loop diagram given in Fig. 8(b) will be given as

$$\frac{d_e}{e}\bigg|_{\chi_k^\pm} = \sum_{i=1,2}\sum_{k=1,2} \mathrm{Im}(\hat{Y}_{H_i^0 e_L e_R^c} C_{\chi_k^+\chi_k^- h})(C_{\chi_k^+\chi_k^-\gamma})^2 \times \frac{\alpha_{em} Q_{\chi_i}^2 Q_e}{16\pi^3 m_{\chi_k^\pm}}(f(z)-g(z)),$$

and the EDM of the neutron for the loop diagram given in Fig. 8(b) will be given as

$$\frac{d_{u/n}}{e}\bigg|_{\chi_k^\pm} = \sum_{i=1,2}\sum_{k=1,2} \mathrm{Im}(\hat{Y}_{H_i^0 e_L e_R^c} C_{\chi_k^+\chi_k^- h})(C_{\chi_k^+\chi_k^-\gamma})^2 \times \frac{\alpha_{em} Q_{\chi_i}^2 Q_u}{16\pi^3 m_{\chi_k^\pm}}(f(z)-g(z)),$$

where $z=\frac{m^2_{\chi_k^\pm}}{m^2_{H_i^0}}$; $f(\frac{m^2_{\chi_1^\pm}}{m^2_{H_1^0}}) - g(\frac{m^2_{\chi_1^\pm}}{m^2_{H_1^0}}) = 10$; $f(\frac{m^2_{\chi_1^\pm}}{m^2_{H_2^0}}) - g(\frac{m^2_{\chi_1^\pm}}{m^2_{H_2^0}}) = 1$, $f(\frac{m^2_{\chi_2^\pm}}{m^2_{H_1^0}}) - g(\frac{m^2_{\chi_2^\pm}}{m^2_{H_1^0}}) = 10$; $f(\frac{m^2_{\chi_2^\pm}}{m^2_{H_2^0}}) - g(\frac{m^2_{\chi_2^\pm}}{m^2_{H_2^0}}) = 0.1$. Considering $(\phi_{\chi_i}-\phi_{y_e})=(\phi_{\chi_i}-\phi_{\tilde{\lambda}_1^0})\sim(0,\frac{\pi}{2}]$, for $m_{\chi_1^\pm}=\mathcal{V}^{\frac{59}{72}}m_{\frac{3}{2}}$, $m_{\chi_2^\pm}=\mathcal{V}^{\frac{2}{3}}m_{\frac{3}{2}}$, $\hat{Y}_{H_i^0 e_L e_R^c}=\mathcal{V}^{-\frac{47}{45}}$, $\hat{Y}_{H_i^0 e_L e_R^c}=\mathcal{V}^{-\frac{17}{18}}$, and the value of the EDM of the electron and neutron will be given as

$$\frac{d_e}{e}\bigg|_{\chi_i} = \frac{d_n}{e}\bigg|_{\chi_i} \sim \frac{\tilde{f}^2\mathcal{V}^{-\frac{8}{3}}}{m_{\frac{3}{2}}} \times 10^{-5}\times 10^{1} \equiv 10^{-33}\ \mathrm{GeV}^{-1} \equiv 10^{-47}\ \mathrm{cm}. \tag{155}$$

3. R-parity violating contribution

For the loop diagram given in Fig. 8(c), the Lagrangian governing the interaction of the neutrino will correspond to R-parity violating interactions given as

$$\mathcal{L} \supset \tilde{\lambda}_{\tilde{\nu}_L u_L u_R^c} \nu_{iL} u_{kL} u_{kR}^c + \tilde{\lambda}_{\tilde{\nu}_L e_L e_R^c} \nu_{iL} e_{jL} e_{jR}^c + \mathrm{H.c.} \tag{156}$$

The contribution of the R-parity violating interaction terms $\hat{\lambda}_{ikk}$ and $\hat{\lambda}^*_{ijj}$ are given by expanding $\mathcal{D}_{\mathcal{A}_1}\mathcal{D}_{\mathcal{A}_3}W$ and $\mathcal{D}_{\mathcal{A}_2}\mathcal{D}_{\mathcal{A}_4}W$ in the fluctuations linear in $\mathcal{A}_1$ around its stabilized VEV. The values of the same have already been calculated in the context of $\mathcal{N}=1$ gauged supergravity action and given as follows:

$$\tilde{\lambda}_{\tilde{\nu}_L e_L e_R^c} \equiv \mathcal{V}^{-\frac{5}{3}} e^{i\phi_{\lambda_e}}, \quad \tilde{\lambda}_{\tilde{\nu}_L u_L u_R^c} \equiv \mathcal{V}^{-\frac{5}{3}} e^{i\phi_{\lambda_u}}, \tag{157}$$

where $e^{i\phi_{\lambda_e}}$ and $e^{i\phi_{\lambda_u}}$ are nonzero phases corresponding to the above-mentioned complex R-parity violating couplings. The EDM of the electron in this case will be

$$\frac{d}{e}\bigg|_{\mathrm{RPV}} = \mathrm{Im}(\tilde{\lambda}_{\tilde{\nu}_L e_L e_R^c}\tilde{\lambda}_{\tilde{\nu}_L u_L u_R^c})\frac{\alpha_{em}Q_u^2 Q_e}{16\pi^3 m_{ej}}(f(z_1)-g(z_1)), \tag{158}$$

and the EDM of the neutron will be given as

$$\frac{d}{u/n}\bigg|_{\mathrm{RPV}} = \mathrm{Im}(\tilde{\lambda}_{\tilde{\nu}_L e_L e_R^c}\tilde{\lambda}_{\tilde{\nu}_L u_L u_R^c})\frac{\alpha_{em}Q_e^2 Q_u}{16\pi^3 m_{ej}}(f(z_2)-g(z_2)), \tag{159}$$

where $z_1=\frac{m_e^2}{m^2_{\nu_{iL}}}$; $z_2=\frac{m_u^2}{m^2_{\nu_{iL}}}$. Using the value of masses $m_{\nu_{iL}}=\mathcal{V}^{\frac{1}{2}}m_{\frac{3}{2}}$, $m_e=0.5\,\mathrm{GeV}$, and $m_u=\mathcal{O}(1)$, $f(\frac{m_e^2/m_u^2}{m^2_{\nu_{iL}}}) = g(\frac{m_e^2/m_u^2}{m^2_{\nu_{iL}}}) = 10^{-27}$, and assuming $(\phi_{\lambda_e}-\phi_{\lambda_u})=(0,\frac{\pi}{2}]$, Eqs. (158) and (159) reduce to give the EDM result as follows:

$$\frac{d}{e}\bigg|_{\mathrm{RPV}} = \frac{d}{n}\bigg|_{\mathrm{RPV}} \sim \mathcal{V}^{-\frac{10}{3}}\times 10^{-2}\times 10^{-27} \equiv 10^{-55}\ \mathrm{GeV}^{-1} \equiv 10^{-70}\ \mathrm{cm}. \tag{160}$$

B. Two-loop-level Barr-Zee Feynman diagrams involving an internal sfermion loop

In this subsection, we evaluate the contribution of a heavy sfermion loop generated by trilinear scalar interactions including Higgs. The loop diagrams are mediated by γh exchange. Unlike one-loop diagrams, here we do not have to consider the mixing of sleptons (squarks) because of the fact that the nonzero phase associated with the complex scalar trilinear interaction is sufficient to generate the nonzero EDM of the elctron/neutron. We first evaluate the contribution of the relevant vertices in the context of $\mathcal{N}=1$ gauged supergravity for the two-loop diagrams shown in Fig. 9.

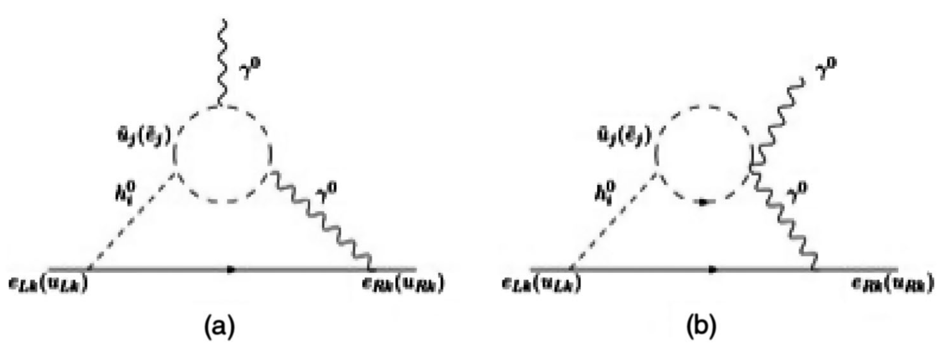

FIG. 9. Two-loop diagrams involving sfermions in the internal loop.

Slepton($\tilde{e}_{jR}$)-slepton($\tilde{e}_{jR}$)-Higgs vertex.—By expanding effective supergravity potential $V|_{\rm EW} \sim e^K K^{\tau_S \bar{\tau}_B} D_{\tau_S} W D_{\bar{\tau}_B} \bar{W} + e^K K^{\tau_B \bar{\tau}_B} |D_{\tau_B} W|^2$ in the fluctuations linear in $\mathcal{Z}_i \to \mathcal{Z}_i + \mathcal{V}^{\frac{1}{36}} M_P$, $\mathcal{A}_3 \to \mathcal{A}_3 + \mathcal{V}^{-\frac{13}{18}} M_P$, the contribution of the term quadratic in $\mathcal{A}_3$ as well as $\mathcal{Z}_i$ is of the order $\mathcal{V}^{-\frac{59}{36}} \langle \mathcal{Z}_i \rangle$, which after giving a VEV to one of the $\mathcal{Z}_i$'s will be given as

$$C_{\tilde{e}_R \tilde{e}_R^* H_i^0} \equiv \frac{1}{\sqrt{(\hat{K}_{\mathcal{Z}_i \bar{\mathcal{Z}}_i})^2 (\hat{K}_{\mathcal{A}_3 \bar{\mathcal{A}}_3})^2}} [\mathcal{V}^{-\frac{59}{36}} \langle \mathcal{Z}_i \rangle] \equiv (\mathcal{V}^{-2} M_P) e^{i\phi_{\tilde{e}_R}}, \tag{161}$$

where $\phi_{\tilde{e}_R}$ is the nonzero phase corresponding to the aforementioned complex scalar three-point interaction vertex. Using the similar procedure, the coefficient of the slepton($\tilde{e}_{jL}$)-slepton($\tilde{e}_{jL}$)-Higgs vertex will be given as

$$C_{\tilde{e}_L \tilde{e}_L^* H_i^0} \equiv \frac{1}{\sqrt{(\hat{K}_{\mathcal{Z}_i \bar{\mathcal{Z}}_i})^2 (\hat{K}_{\mathcal{A}_1 \bar{\mathcal{A}}_1})^2}} [\mathcal{V}^{-\frac{95}{36}} \langle \mathcal{Z}_i \rangle] \equiv (\mathcal{V}^{-\frac{12}{5}} M_P) e^{i\phi_{\tilde{e}_L}}. \tag{162}$$

$\phi_{\tilde{e}_L}$ is the nonzero phase corresponding to this particular complex scalar three-point interaction vertex.

Squark($\tilde{u}_{jR}$)-squark($\tilde{u}_{jR}$)-Higgs vertex.—By expanding $V|_{\rm EW}$ in the fluctuations around $\mathcal{Z}_i \to \mathcal{Z}_i + \mathcal{V}^{\frac{1}{36}} M_P$, $\mathcal{A}_4 \to \mathcal{A}_4 + \mathcal{V}^{-\frac{11}{9}} M_P$, the contribution of the term quadratic in $\mathcal{A}_4$ as well as $\mathcal{Z}_i$ is of the order $\mathcal{V}^{-\frac{23}{36}} \langle \mathcal{Z}_i \rangle$, which after giving the VEV to one of the $\mathcal{Z}_i$, will be given as

$$C_{\tilde{u}_R \tilde{u}_R^* H_i^0} \equiv \frac{1}{\sqrt{(\hat{K}_{\mathcal{Z}_i \bar{\mathcal{Z}}_i})^2 (\hat{K}_{\mathcal{A}_4 \bar{\mathcal{A}}_4})^2}} [\mathcal{V}^{-\frac{23}{36}} \langle \mathcal{Z}_i \rangle] \equiv (\mathcal{V}^{-2} M_P) e^{i\phi_{\tilde{u}_R}}, \tag{163}$$

where $\phi_{\tilde{u}_R}$ is the nonzero phase corresponding to the aforementioned complex scalar three-point interaction vertex.

Squark($\tilde{u}_{jL}$)-squark($\tilde{u}_{jL}$)-Higgs vertex.—By expanding $V|_{\rm EW}$ in the fluctuations around $\mathcal{Z}_i \to \mathcal{Z}_i + \mathcal{V}^{\frac{1}{36}} M_P$, $\mathcal{A}_2 \to \mathcal{A}_2 + \mathcal{V}^{-\frac{1}{3}} M_P$, the contribution of the term quadratic in $\mathcal{A}_2$ as well as $\mathcal{Z}_i$ is of the order $\mathcal{V}^{-\frac{131}{36}} \langle \mathcal{Z}_i \rangle$, which after giving the VEV to one of the $\mathcal{Z}_i$'s will be given as

$$C_{\tilde{u}_L \tilde{u}_L^* H_i^0} \equiv \frac{1}{\sqrt{(\hat{K}_{\mathcal{Z}_i \bar{\mathcal{Z}}_i})^2 (\hat{K}_{\mathcal{A}_2 \bar{\mathcal{A}}_2})^2}} [\mathcal{V}^{-\frac{131}{36}} \langle \mathcal{Z}_i \rangle] \equiv (\mathcal{V}^{-\frac{20}{9}} M_P) e^{i\phi_{\tilde{u}_L}}, \tag{164}$$

where $\phi_{\tilde{u}_L}$ is the nonzero phase corresponding to the aforementioned complex scalar three-point interaction vertex.

The contribution of the slepton($\tilde{e}_{jR}$)-slepton($\tilde{e}_{jR}$)-photon (γ)-photon(γ) vertex will be given by $\bar{\partial}_{\bar{\mathcal{A}}_3} \partial_{\mathcal{A}_3} G_{T_B \bar{T}_B} \times X^{T_B} X^{\bar{T}_B} A^\mu A_\nu$. On solving $\bar{\partial}_{\bar{\mathcal{A}}_3} \partial_{\mathcal{A}_3} G_{T_B \bar{T}_B} \sim \mathcal{V}^{\frac{1}{9}} \mathcal{A}_1^* \mathcal{A}_1$, incorporating the values of X^B as mentioned earlier, the real physical slepton($\tilde{e}_{jR}$)-slepton($\tilde{e}_{jR}$)-photon(γ)-photon(γ) vertex is proportional to

$$C_{\tilde{e}_R \tilde{e}_R^* \gamma\gamma} \equiv \frac{\mathcal{V}^{\frac{1}{9}} \tilde{f}^2 \mathcal{V}^{-\frac{4}{3}}}{\sqrt{(K_{\mathcal{A}_3 \mathcal{A}_3})^2}} \equiv \tilde{f}^2 \mathcal{V}^{-\frac{13}{5}}. \tag{165}$$

The coefficient of the real physical $\tilde{e}_{jL} - \tilde{e}_{jL} \gamma - \gamma$ vertex has been obtained in [38]. The value of the same is given by $C_{\tilde{e}_L \tilde{e}_L^* \gamma\gamma} \equiv \tilde{f}^2 \mathcal{V}^{-3}$. Similarly, the coefficient of the real physical $\tilde{u}_{jR} - \tilde{u}_{jR} - \gamma - \gamma$ vertex will be given by $\bar{\partial}_{\bar{\mathcal{A}}_4} \partial_{\mathcal{A}_4} G_{T_B \bar{T}_B} X^{T_B} X^{\bar{T}_B} A^\mu A_\nu$. On solving, the volume-suppression factor corresponding to this vertex will be given as

$$C_{\tilde{u}_R \tilde{u}_R^* \gamma\gamma} \sim \frac{\text{coefficient of} \, \bar{\partial}_{\bar{\mathcal{A}}_4} \partial_{\mathcal{A}_4} G_{T_B \bar{T}_B}}{\sqrt{(K_{\mathcal{A}_4 \mathcal{A}_4})^2}} \equiv \tilde{f}^2 \mathcal{V}^{-\frac{118}{45}}. \tag{166}$$

The coefficient of the real physical $(\tilde{u}_{jL} - \tilde{u}_{jL} - \gamma - \gamma)$ vertex will be given by $\bar{\partial}_{\bar{\mathcal{A}}_2} \partial_{\mathcal{A}_2} G_{T_B \bar{T}_B} X^{T_B} X^{\bar{T}_B} A^\mu A_\nu$. On solving, the volume-suppression factor corresponding to this vertex will be given as

$$C_{\tilde{u}_L \tilde{u}_L^* \gamma\gamma} \sim \frac{\text{coefficient of} \, \bar{\partial}_{\bar{\mathcal{A}}_2} \partial_{\mathcal{A}_2} G_{T_B \bar{T}_B}}{\sqrt{(K_{\mathcal{A}_2 \mathcal{A}_2})^2}} \equiv \tilde{f}^2 \mathcal{V}^{-\frac{127}{45}}. \tag{167}$$

The contribution of the real scalar-scalar-photon vertices have already been obtained in Sec. II and given as

$$C_{\tilde{e}_L \tilde{e}_L^* \gamma} \equiv (\tilde{f} \mathcal{V}^{\frac{44}{45}}) \tilde{\mathcal{A}}_1 A^\mu \partial_\mu \tilde{\mathcal{A}}_1, \quad C_{\tilde{e}_R \tilde{e}_R^* \gamma} \equiv (\tilde{f} \mathcal{V}^{\frac{53}{45}}) \mathcal{A}_3 A^\mu \partial_\mu \bar{\mathcal{A}}_3,$$
$$C_{\tilde{u}_L \tilde{u}_L^* \gamma} \equiv (\tilde{f} \mathcal{V}^{\frac{53}{45}}) \tilde{\mathcal{A}}_2 A^\mu \partial_\mu \tilde{\mathcal{A}}_2, \quad C_{\tilde{u}_R \tilde{u}_R^* \gamma} \equiv (\tilde{f} \mathcal{V}^{\frac{62}{45}}) \tilde{\mathcal{A}}_4 A^\mu \partial_\mu \tilde{\mathcal{A}}_4. \tag{168}$$

The analytical expression for the EDM involving the sfermion/scalar in an internal loop has been provided in [70]. Using the same, for Fig. 9(a), the EDM of the electron will be given as

$$\frac{d_e}{e}\bigg|_{4.9(a)}^{\rm sfermion} = \sum_{i=1,2} \sum_{j=\tilde{u}_L,\tilde{u}_R} {\rm Im}(\hat{Y}_{H_i^0 e_L e_R^c} C_{H_i^0 j j^*})(C_{jj^*\gamma})^2 \times \frac{\alpha_{em} \eta_c Q_{e_j} q_j^2}{32\pi^3 m_{H_i^0}^2} F(\tilde{z}), \tag{169}$$

and the EDM of the neutron/quark will be given as

$$\left.\frac{d_n}{e}\right|^{\text{sfermion}}_{4.9(a)} = \sum_{i=1,2}\sum_{j=\tilde{e}_L,\tilde{e}_R} \mathrm{Im}(\hat{Y}_{H_i^0 u_L u_R^c} C_{H_i^0 jj^*})(C_{jj^*\gamma})^2 \times \frac{\alpha_{em}\eta_c Q_{u_j} q_j^2}{32\pi^3 m^2_{H_i^0}} F(\tilde{z}), \tag{170}$$

where $z = \frac{m_j^2}{m^2_{H_i^0}}$; $F(z) = -\int_0^1 dx \frac{x(1-x)}{x(1-x)-z}\ln(\frac{x(1-x)}{z})$. Considering $(\phi_{\tilde{u}_{L/R}} - \Phi_{y_e}) = (\phi_{\tilde{e}_{L/R}} - \Phi_{y_u}) = (0, \frac{\pi}{2}]$; $|\hat{Y}_{H_i^0 e_L e_R^c}| \equiv \mathcal{V}^{-\frac{47}{45}}$, $|\hat{Y}_{H_i^0 u_L u_R^c}| \equiv \mathcal{V}^{-\frac{19}{18}}$ and using the value of masses $m_{\tilde{e}_L} = m_{\tilde{e}_R} = m_{\tilde{u}_L} = m_{\tilde{u}_R} = \mathcal{V}^{\frac{1}{2}} m_{\frac{3}{2}}$, $m_{H_1^0} = 125$ GeV, and $m_{H_2^0} = \mathcal{V}^{\frac{59}{72}} m_{\frac{3}{2}}$, we have $F(\frac{m_j^2}{m^2_{H_1^0}}) = 10^{-17}$, $F(\frac{m_j^2}{m^2_{H_2^0}}) = 1$. Incorporating the value of the interaction vertices, Eqs. (169) and (170) reduce to give the EDM results as follows:

$$\left.\frac{d_e}{e}\right|^{\text{sfermion}}_{9(a)} = 10^{-8} \times \mathcal{V}^{-\frac{4}{15}}\tilde{f}^2 \equiv 10^{-15}\ \text{GeV}^{-1} \equiv 10^{-29}\ \text{cm} \quad \text{for } \mathcal{V} = 10^4,$$

$$\left.\frac{d_n}{e}\right|^{\text{sfermion}}_{9(a)} = 10^{-8} \times \mathcal{V}^{-\frac{3}{10}}\tilde{f}^2 \equiv 10^{-15}\ \text{GeV}^{-1} \equiv 10^{-29}\ \text{cm} \quad \text{for } \mathcal{V} = 10^4. \tag{171}$$

For the loop diagram given in Fig. 9(b), the EDM of the electron will be given as

$$\left.\frac{d_e}{e}\right|^{\text{sfermion}}_{9(b)} = \sum_{i=1,2}\sum_{j=\tilde{u}_L,\tilde{u}_R} \mathrm{Im}(\hat{Y}_{H_i^0 e_L e_R^c} C_{H_i^0 jj^*})(C_{jj^*\gamma\gamma}) \times \frac{\alpha_{em} Q_{e_j} q_j^2}{32\pi^3 m^2_{H_i^0}} F(\tilde{z}), \tag{172}$$

and the EDM of the neutron will be given as

$$\left.\frac{d_{n/u}}{e}\right|^{\text{sfermion}}_{9(b)} = \sum_{i=1,2}\sum_{j=\tilde{e}_L,\tilde{e}_R} \mathrm{Im}(\hat{Y}_{H_i^0 u_L u_R^c} C_{H_i^0 jj^*})(C_{jj^*\gamma\gamma}) \times \frac{\alpha_{em} Q_{u_j} q_j^2}{32\pi^3 m^2_{H_i^0}} F(\tilde{z}), \tag{173}$$

where $F(\frac{m_j^2}{m^2_{H_1^0}}) = 10^{-17}$, $F(\frac{m_j^2}{m^2_{H_1^0}}) = 1$. Incorporating the value of the masses and the estimate of the relevant coupling veretx, the EDM of the electron will be

$$\left.\frac{d_e}{e}\right|^{\text{sfermion}}_{9(b)} = 10^{-9} \times \mathcal{V}^{-\frac{17}{3}}\tilde{f}^2 \equiv 10^{-43}\ \text{GeV}^{-1} \equiv 10^{-57}\ \text{cm}. \tag{174}$$

The EDM of the neutron in this case will be given as

$$\left.\frac{d_n}{e}\right|^{\text{sfermion}}_{9(b)} = 10^{-9} \times \mathcal{V}^{-\frac{91}{18}}\tilde{f}^2 \equiv 10^{-40}\ \text{GeV}^{-1} \equiv 10^{-54}\ \text{cm}. \tag{175}$$

The overall contribution of the EDM of the electron as well as the neutron corresponding to the two-loop diagram involving sfermions is

$$\left.\frac{d_{e/n}}{e}\right|^{\text{sfermion}} = \left.\frac{d_{e/n}}{e}\right|^{\text{sfermion}}_{9(a)} + \left.\frac{d_{e/n}}{e}\right|^{\text{sfermion}}_{9(b)} \equiv 10^{-29}\ \text{cm}. \tag{176}$$

C. Two-loop-level Barr-Zee Feynman diagram involving a $W^{\pm}$ boson in the internal loop

In this subsection, we discuss the important contribution of the Barr-Zee diagram involving a W boson as an internal loop. In the one-loop as well as two-loop diagrams discussed so far, we have discussed the contribution mediated by Higgs exchange. The nonzero phases in the one-loop diagram are affected by considering a mixing between the Higgs doublet in a μ-split SUSY model, while in the two-loop diagrams, the phases are affected through a complex effective Yukawa coupling. It has been found in [71] that two-loop graphs involving a W-boson loop can induce an electric dipole moment of d_e of the order of the experimental bound (10^{-27} cm) in the multi-Higgs models provided there is an exchange of Higgs in the Higgs propagator and the CP violation in the neutral Higgs sector is fairly maximal. The approach was given by Weinberg in [72,73]. In these papers, he pointed out that dimension-six purely gluonic operator gives a large value for the EDM of the neutron, which is just below the present experimental bound if one considers CP violation through the exchange of Higgs particles, whose interactions involve one or more complex phases. The approach was extended by Barr and Zee who have found that Higgs exchange can also give an electric dipole moment to the electron of the order of the experimental limits by considering an EDM operator involving a top quark also. In this spirit, we present an analysis of the EDM of the electron/neutron involving a W-boson loop in the context of a μ-split SUSY model, which, as already discussed. involves a light Higgs and a heavy Higgs doublet.

In the notations of Weinberg, the CP-violating phase can appear from the neutral Higgs-boson exchange through imaginary terms in the amplitude, and Higgs propagators are represented as $A(q^2) = \sqrt{2}G_f \sum_n \frac{Z_n}{q^2 + m^2_{H_n}}$, where Z_n is a nonzero phase appearing due to the exchange between the Higgs doublet in the propagator. We address this argument of the generation of the nonzero phase in the $\mathcal{N} = 1$ gauged supergravity action. We first provide the analysis of the

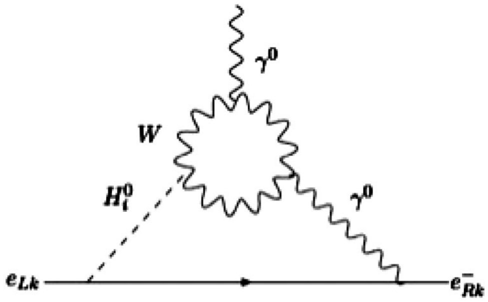

FIG. 10. Two-loop diagram involving a W boson in the internal loop.

required SM-like coupling involved in Fig. 10 in the context of the $\mathcal{N}=1$ gauged supergravity action. The contribution of the W^+-photon-W^- vertex is evaluated by the CP-even interaction term given as [74] $\mathcal{L} = -\mathrm{Re}(f)A^\mu W^-_{\mu\nu}W^{+\nu} + \mathrm{Re}(f)W^{+\mu}W^{-\nu}F_{\mu\nu}$, where $W^-_{\mu\nu} = \partial_\mu W^-_\nu - \partial_\nu W^-_\mu$ and $F_{\mu\nu} = \partial_\mu A_\nu - \partial_\nu A_\mu$, and $\mathrm{Re}(f)$ is a gauge kinetic function, which in our setup is given by the big divisor volume modulus $\mathrm{Re}(T_B) \sim \mathcal{V}^{\frac{1}{18}} \equiv \mathcal{O}(1)$ for Calabi-Yau $\mathcal{V} = 10^5$. Therefore, the volume-suppression corresponding to this interaction vertex is $C_{W^+W^-\gamma} \equiv \mathcal{V}^{\frac{1}{18}} \equiv \mathcal{O}(1)$.

The effective W^+-Higgs-W^- vertex can be evaluated in the effective supergravity action as follows. Consider the gauge kinetic term $\mathrm{Re}(T)F^2$ and then choose the term $C_{1\bar{3}}a_1a_{\bar{3}}$ in $\mathrm{Re}(T_B)$ with the understanding that one first gives the VEV to the predominantly $SU(2)_L$-doublet valued a_1, and then one picks out the $\mathcal{Z}$-dependent contribution in a_3 and also uses the value of the intersection component $C_{1\bar{3}}$. One will, therefore, consider $C_{1\bar{3}}\langle a_1\rangle \mathcal{V}^{-\frac{7}{5}}\mathcal{Z}\frac{p_1\cdot p_2}{\sqrt{K_{\mathcal{A}_1\bar{\mathcal{A}}_1}}\sqrt{K_{\mathcal{Z}\bar{\mathcal{Z}}}}(\sqrt{\mathrm{Re}(T)})^2}$; $K_{\mathcal{Z}\bar{\mathcal{Z}}}|_{M_s} \sim 10^{-5}$, $K_{\mathcal{A}_1\bar{\mathcal{A}}_1}|_{M_s} \sim 10^4$, which at the EW scale we will assume to be $\frac{10^{-5}}{(\mathcal{O}(1))^2}$ and $\frac{10^4}{(\mathcal{O}(1))^2}$. For nonrelativistic gauge bosons, $p_1\cdot p_2 \sim M^2_{W/Z}$, $\mathrm{Re}(T)|_{\mathrm{EW}} \sim \mathcal{O}(1)M_P \sim v\mathcal{V}^3$ GeV, $C_{1\bar{3}} \sim \mathcal{V}^{\frac{29}{18}}$, $\langle a_1\rangle|_{\mathrm{EW}} \sim \mathcal{O}(1)\mathcal{V}^{-\frac{2}{9}}M_P$ [related to the requirement of obtaining $\mathcal{O}(10^2)$ GeV W/Z-boson mass at the EW scale; see [38]]. We, thus, obtain the following: $\mathcal{V}^{\frac{29}{18}} \times (\mathcal{O}(1))^2 \times \mathcal{O}(1) \times \mathcal{V}^{-\frac{2}{9}} \times \mathcal{V}^{-\frac{7}{5}}M^2_{W/Z} \times \frac{\sqrt{10}}{(\mathcal{O}(1)\times v \times \mathcal{V}^3)} \sim (\mathcal{O}(1))^2 \times \sqrt{10}\mathcal{V}^{-3}\frac{M^2_{W/Z\,\mathrm{in\,GeV}}}{v(\mathrm{GeV})}$. Now, in the superspace notation, the kinetic terms for the gauge field are generically written as $\int d^2\theta f_{ab}(\Phi)W^aW^b$ where W^a is the gauge-invariant superfield strength and $W = W^aT^a$ for a non-Abelian group—as f_{ab} is an *a priori* arbitrary holomorphic function of Φ. Consider, hence, $\Phi = T, f \sim e^T$ and look at $\int d^2\theta (T)^{2m+1}_{\theta,\bar{\theta}=0}W^2$, which will consist of $(\mathcal{O}(1)^2 \times C_{1\bar{3}}\langle a_1\bar{a}_{\bar{3}}\rangle)^{2m} \times \sqrt{10} \times \mathcal{V}^{-3}\frac{M^2_{W/Z\,\mathrm{in\,GeV}}}{\mathrm{GeV}}$, which, e.g., for $m=2$ yields $(\mathcal{O}(1))^2 \times \sqrt{10} \times \mathcal{V}^{\frac{8}{3}-3} \times \frac{M^2_{W/Z\,\mathrm{in\,GeV}}}{v(\mathrm{GeV})}$ or for $\mathcal{V} \sim 10^4$, one obtains $\mathcal{O}(1)\frac{M^2_{W/Z(\mathrm{in\,GeV})}}{v(\mathrm{in\,GeV})}$. Utilizing this, at the EW scale, $C_{W^+H^0_iW^-} \equiv \frac{M^2_W}{v}e^{i\phi_W}$. The value of the complex Yukawa coupling to be used to evaluate the EDM corresponding to Fig. 10 has already been obtained in [38] and given as $\hat{Y}_{H^0_ie_Le^c_R} \sim \mathcal{V}^{-\frac{47}{45}}e^{i\phi_{y_e}}$ and of $\hat{Y}_{H^0_iu_Lu^c_R} \sim \mathcal{V}^{-\frac{19}{18}}e^{i\phi_{y_u}}$. The matrix amplitude as well as the analytical expression for the W boson related the loop diagrams has been worked out in [71]. We utilize the same in a modified form to get the numerical estimate of the EDM corresponding to the loop diagram given in Fig. 10,

$$\frac{d}{e}\bigg|_W = \frac{\alpha}{(4\pi)^3M^2_W}C_{W^+W^-\gamma}\sum_i \mathrm{Im}(\hat{Y}_{H^0_ie_Le^c_R}C_{W^+H^0_iW^-}) \times\left[5g(z^W_i) + 3f(z^W_i) + \frac{3}{4}(g(z^W_i) + h(z^W_i))\right], \tag{177}$$

where $f(z)$ and $g(z)$ are already defined in Sec. IVA and

$$h(z) = \frac{z}{2}\int_0^1 dx\frac{1}{x(1-x)-z} \times\left(\frac{z}{x(1-x)-z}\ln\left(\frac{x(1-x)-z}{x}\right) - 1\right), \tag{178}$$

where $z^W_i = \frac{m^2_{H^0_i}}{m^2_W}$. Considering $(\phi_W - \phi_{y_e}) = (0, \frac{\pi}{2}]$, and using the values $m_{H^0_1} = 125$ GeV and $m_{H^0_2} = \mathcal{V}^{\frac{59}{72}}m_{\frac{3}{2}}$, we get $f(\frac{m^2_{H^0_i}}{m^2_W}) = g(\frac{m^2_{H^0_i}}{m^2_W}) = h(\frac{m^2_{H^0_i}}{m^2_W}) = \mathcal{O}(1)$, and the EDM result for the electron will be given as

$$\frac{d_e}{e}\bigg|_W \sim \frac{\alpha}{(4\pi)^3}\frac{1}{v} \times \mathcal{V}^{-\frac{47}{45}} \equiv 10^{13}\ \mathrm{GeV}^{-1} \equiv 10^{-27}\ \mathrm{cm}. \tag{179}$$

Similarly, by considering $(\phi_W - \phi_{y_u}) = (0, \frac{\pi}{2}]$, the EDM of the neutron will be given as

$$\frac{d_n}{e}\bigg|_W \sim \frac{\alpha}{(4\pi)^3}\frac{1}{v} \times \mathcal{V}^{-\frac{19}{18}} \equiv 10^{13}\ \mathrm{GeV}^{-1} \equiv 10^{-27}\ \mathrm{cm}. \tag{180}$$

D. Two-loop-level rainbow-type contribution to the electric dipole moment

The two-loop-level analysis of the supersymmetric effects to the fermion electric dipole moment has been extended by considering rainbow diagrams in addition to famous Barr-Zee diagrams with the expectation that they might give a significant contribution to fermionic EDM.

The importance of these diagrams is discussed in detail in [68]. In this subsection, we estimate the contribution of two-loop rainbow-type of diagrams involving R-parity conserving supersymmetric interaction vertices and R-parity violating vertices. The CP-violating phases appear from the diagonalized eigenstates in the inner loop as well as from complex effective Yukawa couplings in the Higgsino sector. The Feynman diagrams have been classified based on different types of inner one-loop insertions. One corresponds to the one-loop effective Higgsino-gaugino-gauge boson vertex, and the other corresponds to the one-loop effective Higgsino-gaugino transition. The matrix amplitudes as well as the analytic expressions to estimate the EDM for the above rainbow diagrams are calculated in [75] to the first order in the external momentum carried by the gauge boson. We utilize their expressions to get the order of magnitude of the EDM of the electron as well as the neutron in our case.

1. R-parity conserving rainbow-type contributions

For the loop diagrams given in Figs. 11 and 12, the result of the EDM will be given by the following formulas, respectively,

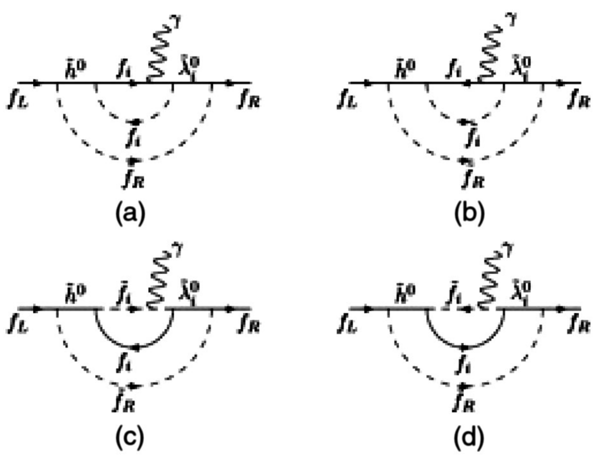

FIG. 11. Two-loop-level rainbow-type diagrams involving the Higgsino-gaugino-gauge boson vertex.

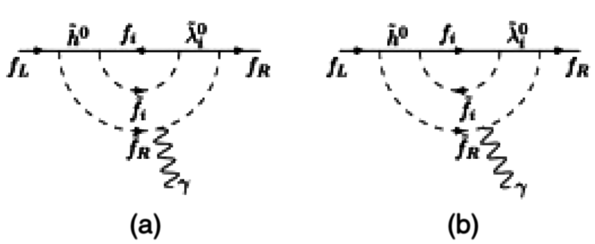

FIG. 12. Two-loop-level rainbow-type diagrams involving Higgsino-gaugino transition.

$$
\begin{aligned}
d_f^1 \approx & \sum \frac{n_c(Q_f+Q'_f)C_{\tilde{h}^0 f_L \tilde{f}_R^*}C_{\tilde{h}^0 f_i f_i^*}}{64\pi^3}\sum_{n=1,2}|m_{\lambda_n^0}|\sin(\delta_f-\theta_n)\\
&\times\frac{e(g'^{(n)}_{\tilde{f}_L}-g'^{(n)}_{\tilde{f}_R})}{4\pi}\sin\theta_f\cos\theta_f\\
&\times\sum_{\tilde{f}=\tilde{f}_L,\tilde{f}_R} s g'^{(n)}_{\tilde{f}}[F'(|m_{\lambda_n}|^2,|\mu|^2,m^2_{\tilde{f}_{L/R}},m^2_{\tilde{f}_2})\\
&-F'(|m_{\lambda_n^0}|^2,|\mu|^2,m^2_{\tilde{f}_{L/R}},m^2_{\tilde{f}_1})], \qquad (181)
\end{aligned}
$$

$$
\begin{aligned}
d_f^2 = & \sum_f \frac{n_c Q'_{f_R} C_{\tilde{h}^0 f_L \tilde{f}_R^*}C_{\tilde{h}^0 f_i f_i^*}}{64\pi^3}\sum_{n=1,2}|m_{\lambda_n^0}|\sin(\delta_f-\theta_n)\\
&\times\frac{e(g'^{(n)}_{\tilde{f}_L}+g'^{(n)}_{\tilde{f}_R})}{4\pi}\sin\theta_f\cos\theta_f\\
&\times\sum_{\tilde{f}=\tilde{f}_L,\tilde{f}_R} g'^{(n)}_{\tilde{f}} m^2_{\tilde{F}}[F''(|m_{\lambda_n}|^2,|\mu|^2,m^2_{\tilde{f}_{L/R}},,m^2_{\tilde{f}_1})\\
&-F''(|m_{\lambda_n}|^2,|\mu|^2,m^2_{\tilde{f}_{L/R}},,m^2_{\tilde{f}_2})], \qquad (182)
\end{aligned}
$$

where $n_c = 3$ for the inner quark-squark loop and $n_c = 1$ for the inner lepton-slepton loop. The fields $\tilde{f}_1$ and $\tilde{f}_2$ correspond to the mass eigenstates of the sfermion $\tilde{f}$. The value of constant s is $+1$ for the left-handed sfermion $\tilde{f}_L$ and -1 for the right-handed sfermion $\tilde{f}_R$. The effective electric charges are given by $Q'_f = C_{f_i f_i^* \gamma} Q_f$ and $Q'_{f_R} = C_{f_R f_R^* \gamma}$. The interaction vertices $C_{\tilde{h}^0 f_L f_R^*}$ and $C_{\tilde{h}^0 f f^*}$ correspond to effective Yukawa couplings. $g'^{(n)}_{\tilde{f}_L}$ and $g'^{(n)}_{\tilde{f}_R}$ denote the effective gauge couplings corresponding to supersymmetric sfermions, and the functions F' and F'' are defined in [75]. The effective Yukawa's as well as the gauge interaction vertices are already calculated in Sec. IV. The magnitude of the values of the same are

$$
\begin{aligned}
&|C^{\tilde{h}^0 e_L \tilde{e}_R^*}| \equiv \mathcal{V}^{-\frac{9}{5}}, \qquad |C^{\tilde{h}^0 e_i \tilde{e}_i{}^*}| \equiv \mathcal{V}^{-\frac{10}{9}}, \qquad |C^{\tilde{e}_i \tilde{e}_i \gamma}| \equiv \tilde{f}\mathcal{V}^{\frac{53}{45}},\\
&|C^{\tilde{h}^0 u_L \tilde{u}_R^*}| \equiv \mathcal{V}^{-\frac{5}{3}}, \qquad |C^{\tilde{h}^0 u_i \tilde{u}_i{}^*}| \equiv \mathcal{V}^{-\frac{10}{9}},\\
&|C^{\tilde{u}_i \tilde{e}_i \gamma}| \equiv \tilde{f}\mathcal{V}^{\frac{53}{45}}, \qquad g'^{(n)}_{\tilde{e}_R} \equiv |C^{e_R \tilde{e}_R^* \lambda_i^0}| \equiv \tilde{f}\mathcal{V}^{-\frac{3}{5}},\\
&g'^{(n)}_{\tilde{e}} \equiv |C^{e_i \tilde{e}_i^* \lambda_i^0}| \equiv \tilde{f}\mathcal{V}^{-\frac{3}{5}}, g'^{(n)}_{\tilde{e}_R} \equiv |C^{u_R \tilde{u}_R^* \lambda_i^0}| \equiv \tilde{f}\mathcal{V}^{-\frac{3}{5}},\\
&g'^{(n)}_{\tilde{u}} \equiv |C^{u_i \tilde{u}_i^* \lambda_i^0}| \equiv \tilde{f}\mathcal{V}^{-\frac{3}{5}}; \quad i = 1,2. \qquad (183)
\end{aligned}
$$

Using the value of $m_{\lambda_1^0} = m_{\lambda_2^0} = \mathcal{V}^{\frac{2}{3}} m_{\frac{3}{2}}$, $m^2_{\tilde{e}_1} = m^2_{\tilde{u}_1} = \mathcal{V} m^2_{\frac{3}{2}} + m^2_{\tilde{e}_{12}}$, $m^2_{\tilde{e}_2} = m^2_{\tilde{u}_2} = \mathcal{V} m^2_{\frac{3}{2}} - m^2_{\tilde{e}_{12}}$, $m_{\tilde{h}^0} \equiv \mathcal{V}^{\frac{59}{72}} m_{\frac{3}{2}}$, we get

$$
\begin{aligned}
&F'(|m_{\lambda_n}|^2,|\mu|^2,m^2_{\tilde{f}_{L/R}},m^2_{\tilde{f}_2}) - F'(|m_{\lambda_n}|^2,|\mu|^2,m^2_{\tilde{f}_{L/R}},m^2_{\tilde{f}_1})\\
&\quad\equiv 10^{-24},\\
&F'(|m_{\lambda_n}|^2,|\mu|^2,m^2_{\tilde{f}_{L/R}},m^2_{\tilde{f}_2}) - F'(|m_{\lambda_n}|^2,|\mu|^2,m^2_{\tilde{f}_{L/R}},m^2_{\tilde{f}_1})\\
&\quad\equiv 10^{-43}. \qquad (184)
\end{aligned}
$$

Incorporating the above results in the analytical expression as given in (181) and (182) [with the assumption that the value of the phase factor associated with all effective R-parity conserving Yukawa couplings are of $\mathcal{O}(1)$],

$$d_e^1/e \equiv \tilde{f}^3\mathcal{V}^{-\frac{18}{5}} \times \mathcal{V}^{\frac{2}{3}}m_{\frac{3}{2}} \times (10^{-24})\ \mathrm{GeV}^{-2} \equiv 10^{-57}\ \mathrm{cm},$$
$$d_e^2/e \equiv \tilde{f}^3\mathcal{V}^{-\frac{18}{5}} \times \mathcal{V}^{\frac{5}{3}}m_{\frac{3}{2}}^3 \times (10^{-43})\ \mathrm{GeV}^{-4} \equiv 10^{-55}\ \mathrm{cm}, \tag{185}$$

and similarly,

$$d_u^1/e \equiv \tilde{f}^3\mathcal{V}^{-\frac{10}{3}} \times \mathcal{V}^{\frac{2}{3}}m_{\frac{3}{2}} \times (10^{-24})\ \mathrm{GeV}^{-2} \equiv 10^{-56}\ \mathrm{cm},$$
$$d_u^2/e \equiv \tilde{f}^3\mathcal{V}^{-\frac{10}{3}} \times \mathcal{V}^{\frac{5}{3}}m_{\frac{3}{2}}^3 \times (10^{-43})\ \mathrm{GeV}^{-4} \equiv 10^{-54}\ \mathrm{cm}. \tag{186}$$

So, the final EDM of the electron as well as the quark or neutron in the case of R-parity conserving supersymmetric Feynman diagrams are given as

$$d_e/e = d_e^1/e + d_e^2/e \equiv 10^{-55}\ \mathrm{cm},$$
$$d_u/e = d_u^1/e + d_u^2/e \equiv 10^{-54}\ \mathrm{cm}. \tag{187}$$

2. R-parity violating rainbow-type contribution

The similar kind of Feynman diagrams can be drawn by replacing the neutral Higgsino component with the Dirac massless neutrino in Figs. 11 and 12. The formulas of the EDM of the fermion f for two types of Feynman diagrams as defined in [76] are given below

$$d_f^1 = \sum_{n=1,2} \mathrm{Im}(C_{\tilde{h}^0 f_L \tilde{f}_R^*} C_{\tilde{h}^0 f_i \tilde{f}_i^*} e^{i(\theta_n - \delta_{f_j})}) \frac{(Q_f + Q'_f)n_c}{64\pi^3}|m_{\lambda_n^0}| \times \frac{e(g^{(n)}_{\tilde{f}_L} - g^{(n)}_{\tilde{f}_R})}{4\pi}\sin\theta_{f_j}\cos\theta_{f_j} \times \sum_{\tilde{f}=\tilde{f}_L,\tilde{f}_R} sg^{(n)}_{\tilde{f}}[f\prime(|m_{\lambda_n^0}|^2, 0, m^2_{\tilde{f}_{L/R}}, m^2_{\tilde{f}_{1j}}) - f\prime(|m_{\lambda_n}|^2, 0, m^2_{\tilde{f}_{L/R}}, m^2_{\tilde{f}_{2j}})], \tag{188}$$

$$d_f^2 = -\sum_{n=1,2} \mathrm{Im}(C_{\tilde{h}^0 f_L \tilde{f}_R^*} C_{\tilde{h}^0 f_i \tilde{f}_i^*} e^{i(\theta_n - \delta_{f_j})}) \frac{Q'_{f_R} n_c}{64\pi^3}|m_{\lambda_n^0}| \times \frac{e(g^{(n)}_{\tilde{f}_L} + g^{(n)}_{\tilde{f}_R})}{4\pi}\sin\theta_{f_j}\cos\theta_{f_j} \times \sum_{\tilde{f}=\tilde{f}_L,\tilde{f}_R} g^{(n)}_{\tilde{f}} m^2_{\tilde{f}_k}[f''(|m_{\lambda_n^0}|^2, 0, m^2_{\tilde{f}_{L/R}}, m^2_{\tilde{f}_1}) - f''(|m_{\lambda_n}|^2, 0, m^2_{\tilde{f}_{L/R}}, m^2_{\tilde{f}_2})]. \tag{189}$$

The interaction vertices $C_{\nu^0 f_L \tilde{f}_R^*}$ and $C_{\nu^0 f_i \tilde{f}_i^*}$ correspond to effective R-parity violating couplings. $g\prime^{(n)}_{\tilde{f}_L}$ and $g\prime^{(n)}_{\tilde{f}_R}$ denote effective gauge couplings corresponding to supersymmetric sfermions. The functions $F\prime$ and F'' are defined in [76].

Using the value of $m_{\lambda_1^0} = m_{\lambda_2^0} = \mathcal{V}^{\frac{2}{3}}m_{\frac{3}{2}}$, $m^2_{\tilde{e}_1} = m^2_{\tilde{u}_1} = \mathcal{V}m^2_{\frac{3}{2}} + m^2_{\tilde{e}_{12}}$, $m^2_{\tilde{e}_2} = m^2_{\tilde{u}_2} = \mathcal{V}m^2_{\frac{3}{2}} - m^2_{\tilde{e}_{12}}$, we get

$$F\prime(|m_{\lambda_n^0}|^2, 0, m^2_{\tilde{f}_{L/R}}, m^2_{\tilde{f}_2}) - F\prime(|m_{\lambda_n^0}|^2, |\mu|^2, m^2_{\tilde{f}_{L/R}}, m^2_{\tilde{f}_1}) \equiv 10^{-22},$$
$$F\prime(|m_{\lambda_n^0}|^2, 0, m^2_{\tilde{f}_{L/R}}, m^2_{\tilde{f}_2}) - F\prime(|m_{\lambda_n^0}|^2, |\mu|^2, m^2_{\tilde{f}_{L/R}}, m^2_{\tilde{f}_1}) \equiv 10^{-42}. \tag{190}$$

The contribution of R-parity violating vertices are already calculated in [38] in the context of $\mathcal{N} = 1$ gauged supergravity action. The values of the same are as follows:

$$|C^{\nu^0 e_L \tilde{e}_R^*}| \equiv |C^{\nu^0 e_i \tilde{e}_i^*}| \equiv \mathcal{V}^{-\frac{5}{3}},$$
$$|C^{\nu^0 u_L \tilde{u}_R^*}| \equiv |C^{\nu^0 u_i \tilde{u}_i^*}| \equiv \mathcal{V}^{-\frac{5}{3}}; \quad i = 1, 2. \tag{191}$$

Incorporating the values of the above-mentioned R-parity violating interaction vertices and the values of the effective gauge couplings in the analytic expressions given in Eqs. (188) and (189) [with the assumption that the value of the phase factor associated with all effective R-parity violating Yukawa couplings is of $\mathcal{O}(1)$],

TABLE III. Results of the EDM of the electron/neutron for all possible two-loop diagrams.

Two-loop particle exchange	Origin of complex phase	$d_e(e\ \mathrm{cm})$	$d_n(e\ \mathrm{cm})$
$h_i^0\gamma f$	Complex effective Yukawa couplings	10^{-36}	10^{-36}
$h_i^0\gamma\chi_i^{\pm}$	"	10^{-47}	10^{-47}
$\tilde{f}f\gamma$	"	10^{-70}	10^{-70}
$\tilde{f}_i^0 h_i^0\gamma$	"	10^{-29}	10^{-29}
$\gamma W^{\pm}h_i^0$	Higgs exchange	10^{-27}	10^{-27}
$\tilde{h}^0\tilde{f}\lambda_i^0$ (rainbow type)	Diagonalized sfermion mass eigenstates and effective Yukawas	10^{-55}	10^{-54}
$\nu^0\tilde{f}\lambda_i^0$ (rainbow type)	"	10^{-52}	10^{-52}

$$d_e^1/e = d_u^1/e \equiv \tilde{f}^3 \mathcal{V}^{-\frac{10}{3}} \times \mathcal{V}^{\frac{2}{3}} m_{\frac{3}{2}} \times (10^{-22})\,\mathrm{GeV}^{-2} \equiv 10^{-53}\ \mathrm{cm},$$
$$d_e^1/e = d_u^2/e \equiv \tilde{f}^3 \mathcal{V}^{-\frac{10}{3}} \times \mathcal{V}^{\frac{5}{3}} m^3_{\frac{3}{2}} \times (10^{-42})\,\mathrm{GeV}^{-4} \equiv 10^{-52}\ \mathrm{cm}. \quad (192)$$

So, the final EDM of the electron as well as the quark or neutron in the case of R-parity violating Feynman diagrams is given as

$$d_n/e = d_e/e \equiv 10^{-52}\ \mathrm{cm}. \quad (193)$$

The results of all two-loop diagrams contributing to the EDM of the electron/neutron are summarized in Table III.

VI. SUMMARY AND CONCLUSIONS

To summarize, we have performed a quantitative order-of-magnitude analysis of the EDM of the electron and neutron in a phenomenological model which provides a local realization of large volume $D3/D7$ μ-split supersymmetry that could possibly locally be obtained in the framework of four Wilson line moduli living on the world volume of fluxed stacks of space-time filling $D7$-branes wrapped around the big divisor and two position moduli of a mobile space-time filling $D3$-brane restricted to (nearly) a special Lagrangian three cycle of a Swiss-cheese Calabi-Yau. The proposed phenomenological model is governed by a superheavy gaugino and Higgsino mass parameter in addition to heavy sfermion masses except one light Higgs (obtained by considering a linear combination of the eigenstates of the Higgs doublets at the EW scale). Because of the presence of a heavy gaugino/Higgsino mass parameter, one cannot ignore one-loop diagrams mediated by gaugino/Higgsinos and sfermions as compared to two-loop diagrams as traditional split SUSY models do. Keeping this in mind, we have taken into account the complete set of one-loop graphs and the dominant Higgs-mediated Barr-Zee diagrams. The nonzero CP-violating phase corresponding to the dimension-five nonrenormalizable EDM operator can be made to appear at the one-loop and two-loop levels from the off-diagonal component of a scalar mass matrix and complex effective Yukawa couplings, respectively, in the context of $\mathcal{N} = 1$ gauged supergravity action. We have considered the order of phases to exist in $(0,\frac{\pi}{2}]$. We have also shown that for a given choice of VEVs of a Wilson line as well as position moduli, the phases corresponding to effective Yukawa couplings do not change in the renormalization group flow from string scale down to the electroweak scale. The relevant interaction vertices have been calculated in the context of $\mathcal{N} = 1$ effective gauged supergravity action. Having described the aforementioned model, we estimate all possible one-loop as well as two-loop diagrams. In the one-loop graphs involving sfermions, the neutralino-mediated loop diagrams give the dominant contributions to the electron (neutron) EDM values as compared to gaugino-mediated one-loop diagrams and the diagrams involving R-parity violating vertices, because in $\mathcal{N} = 1$ gauged supergravity, gaugino interaction vertices are dependent on suppressed dilute non-Abelian fluxes, and the contribution of R-parity violating vertices are generally suppressed. However, all of the three-loop diagrams give a very suppressed contribution to the electron and neutron EDM. Next, we considered one-loop diagrams involving Higgs and other supersymmetric/SM particles. By considering Standard-Model-like fermions with Higgs in a loop, we get the electron EDM estimate ($d_e/e \equiv 10^{-34}$ cm) and neutron EDM estimate ($d_n/e \equiv 10^{-33}$ cm) considerably higher than the value predicted by the Standard Model. Interestingly, by considering one-loop diagrams involving a chargino and Higgs, the electron EDM value turns out to be ($d_e/e \equiv 10^{-32}$ cm); i.e., one gets a healthy EDM of the electron even in the presence of a superheavy chargino in the loop. All of the above one-loop diagrams involve MSSM-like superfields, and CP-violating phases appear from visible sector fields only. For a full-fledged analysis, we have taken into account a Goldstino supermultiplet also as the physical degrees of freedom in the one-loop diagrams. As the sGoldstino corresponds to the bosonic component of the superfield corresponding to which there is a supersymmetry breaking and the same occurs maximally in our large volume $D3/D7$ model via the complex big divisor volume modulus (τ_B), we have identified the sGoldstino field with a complex τ_B field. Since, the fermionic component Goldstino gets absorbed into the gravitino and becomes a longitudinal component of the massive gravitino, we basically consider one-loop diagrams involving a gravitino and an sGoldstino in the loop. In such kind of loop diagrams, CP-violating phases appear from hidden sector fields. However, by evaluating the matrix amplitudes of these loop diagrams, we get a very suppressed contribution of the electron and neutron EDM. The results of all one-loop diagrams are summarized in Table II. In the case of two-loop diagrams, we have evaluated the contribution of Barr-Zee diagrams involving sfermions/fermions in an internal loop and mediated via γh exchange and an R-parity violating diagram involving fermions but mediated via $\tilde{f} h$ exchange. Here, the two-loop Barr-Zee diagrams involving heavy sfermions and a light Higgs give a most dominant contribution of the EDM ($d_{(e/n)}/e \equiv 10^{-29}$ cm) as compared to two-loop diagrams involving only SM-like particles. With substantial fine-tuning in Calabi-Yau volume, one can hope to produce EDM results the same as experimental limits. Next, inspired by the approach given in [71–73,77] to obtain a large EDM value (almost the same as an experimental bound) from Barr-Zee diagrams involving top quarks and a W-bosons loop in multi-Higgs models, we have provided an estimate of the same using two Higgs doublets given in the context of μ-split SUSY. By showing the possibility of obtaining the numerical estimate of all SM-like vertices relevant for these diagrams to be same as their standard values in the context of the $\mathcal{N}=1$ gauged supergravity model, we have also produced the EDM ($d_{(e,n)}/e \equiv 10^{-27}$ cm) in the case of a Barr-Zee diagram involving a W boson. As evaluated explicitly, we have shown that two-loop rainbow diagrams give a very suppressed contribution as compared to Barr-Zee diagrams. The results of all two-loop diagrams are summarized in Table III. Thus, we conclude that in our large volume $D3/D7$ μ-split SUSY model, despite the presence of very heavy supersymmetric scalars/fermions in the loops, we are able to produce a contribution to the electric dipole moment of both the electron as well as the neutron close to the experimental bound at the two-loop level and a sizable contribution even at the one-loop level.

All of the above results have been obtained in the context of the model which can be constructed locally near a particular nearly special Lagrangian three cycle of a Swiss-cheese Calabi-Yau threefold. It would be interesting to determine the global embedding of our model. Further, in the $D3/D7$ setup described above, we have shown the possibility of identification of Wilson line moduli only with first- or second-generation quarks and leptons. By extending the setup to include Wilson line moduli identifiable with second- and third-generation quarks, one hopes to obtain via the one-loop and two-loop Barr-Zee diagrams involving fermions, the value of the electron/neutron EDM very close to the experimental bound for a given choice of the internal complex threefold volume.

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APPENDIX: CHARGINO MASS MATRIX

The chargino mass matrix is formed by mixing (charged) winos and a Higgsino after electroweak symmetry breaking. In $\mathcal{N}=1$ gauged supergravity, the interaction vertex corresponding to the Higgs-gaugino-Higgsino term is given by $\mathcal{L}=g_{YM}g_{T^B\mathcal{Z}^i}X^B\tilde{H}^i\lambda^i+\partial_{\mathcal{Z}_i}T_BD^B\tilde{H}^i\lambda^i$ where λ^i corresponds to a gaugino (such as the bino/wino). Expanding the same in the fluctuations linear in Z_i, we have

$$g_{YM}g_{T^B\mathcal{Z}^i}X^B=\tilde{f}\mathcal{V}^{-2}\frac{\mathcal{Z}_i}{M_P},\quad (\partial_{\mathcal{Z}_i}T_B)D^B\equiv\tilde{f}\mathcal{V}^{-\frac{4}{3}}\frac{\mathcal{Z}_i}{M_P}. \tag{A1}$$

After giving the VEV to $\mathcal{Z}_i$, the interaction vertex corresponding to mixing between the gaugino and Higgsino will be given as

$$C^{\tilde{\lambda}^--\tilde{H}_d^-}/C^{\tilde{\lambda}^+-\tilde{H}_u^+}=\frac{\tilde{f}\mathcal{V}^{-\frac{4}{3}}}{\sqrt{K_{\mathcal{Z}_i\mathcal{Z}_i}K_{\mathcal{Z}_i\mathcal{Z}_i}}}\equiv\tilde{f}\mathcal{V}^{-\frac{1}{3}}\frac{v}{M_P},$$
$$\text{where } v=246\text{ GeV}. \tag{A2}$$

For Higgsino doublets $\tilde{H}_u=(\tilde{H}_u^0,\tilde{H}_u^+)$, $\tilde{H}_d=(\tilde{H}_d^-,\tilde{H}_d^0)$, the chargino mass matrix is given as

$$M_{\tilde{\chi}^-}=\begin{pmatrix} M^2_{\tilde{H}_d^-} & C^{\tilde{\lambda}^--\tilde{H}_d^-}\\ C^{\tilde{\lambda}^--\tilde{H}_d^-} & M^2_{\tilde{\lambda}^-}\end{pmatrix},$$
$$M_{\tilde{\chi}^+}=\begin{pmatrix} M^2_{\tilde{H}_u^+} & C^{\tilde{\lambda}^+-\tilde{H}_u^+}\\ C^{\tilde{\lambda}^+-\tilde{H}_u^+} & M^2_{\tilde{\lambda}^+}\end{pmatrix}. \tag{A3}$$

Incorporating the values of $M_{\tilde{\lambda}^+}=M_{\tilde{\lambda}^-}=\mathcal{V}^{\frac{2}{3}}m_{\frac{3}{2}}$, $M_{\tilde{H}_d^-}=M_{\tilde{H}_u^+}=\mathcal{V}^{\frac{59}{72}}m_{\frac{3}{2}}$, and $m_{\frac{3}{2}}=\mathcal{V}^{-2}M_P$ at the electroweak scale, we have

$$M_{\tilde{\chi}^\pm}=\begin{pmatrix}\mathcal{V}^{-\frac{4}{3}} & \frac{v}{M_P}\tilde{f}\mathcal{V}^{-\frac{1}{3}}\\ \frac{v}{M_P}\tilde{f}\mathcal{V}^{-\frac{1}{3}} & \mathcal{V}^{-\frac{85}{72}}\end{pmatrix} \tag{A4}$$

giving eigenvalues

$$\left\{\frac{M_P^2\mathcal{V}^{4/3}+M_P^2\mathcal{V}^{85/72}-\sqrt{M_P^4\mathcal{V}^{8/3}-2M_P^4\mathcal{V}^{181/72}+M_P^4\mathcal{V}^{85/36}+4\tilde{f}^2M_P^2v^2\mathcal{V}^{109/36}}}{2M_P^2\mathcal{V}^{157/72}},\right.$$
$$\left.\frac{M_P^2\mathcal{V}^{4/3}+M_P^2\mathcal{V}^{85/72}+\sqrt{M_P^4\mathcal{V}^{8/3}-2M_P^4\mathcal{V}^{181/72}+M_P^4\mathcal{V}^{85/36}+4\tilde{f}^2M_P^2v^2\mathcal{V}^{109/36}}}{2M_P^2\mathcal{V}^{157/72}}\right\}M_P, \tag{A5}$$

and normalized eigenvectors

$$\tilde{\chi}_1^+ = -\tilde{H}_u^+ + \left(\frac{v}{M_P}\tilde{f}\mathcal{V}^{\frac{5}{6}}\right)\tilde{\lambda}_i^+, \qquad \tilde{\chi}_1^- = -\tilde{H}_d^- + \left(\frac{v}{M_P}\tilde{f}\mathcal{V}^{\frac{5}{6}}\right)\tilde{\lambda}_i^-, \quad \text{and} \quad m_{\tilde{\chi}_1^\pm} \equiv \mathcal{V}^{\frac{59}{72}}m_{\frac{3}{2}},$$

$$\tilde{\chi}_2^+ = \tilde{\lambda}_i^+ + \left(\frac{v}{M_P}\tilde{f}\mathcal{V}^{\frac{5}{6}}\right)\tilde{H}_u^+, \qquad \tilde{\chi}_2^- = \tilde{\lambda}_i^- + \left(\frac{v}{M_P}\tilde{f}\mathcal{V}^{\frac{5}{6}}\right)\tilde{H}_d^-, \quad \text{and} \quad m_{\tilde{\chi}_2^\pm} \equiv \mathcal{V}^{\frac{2}{3}}m_{\frac{3}{2}}. \tag{A6}$$

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