strengths [4]. Also, the inflaton coupling to fermionic partners of $\chi_{1,2}$ follows naturally from SUSY. The prospects for fermionic preheating will thus be the same as those for the bosonic case.

B. Supersymmetric Flat Directions

A key property of SUSY theories is the presence of flat directions in field space along which the potential identically vanishes (in the limit of unbroken SUSY). Such scalar fields (which are complex) can therefore obtain large VEVs along these special directions at no energy cost. These flat directions, which can be interpreted as a degeneracy of the vacuum state of SUSY theories, arise because of cancellations between fields of opposite charges in the D-term potential. A powerful tool for finding the flat directions has been developed in [62–64] (for reviews see [65**?** , 66]). Flat directions are classified by gauge-invariant monomials $\prod_{i=1}^{n} X_i$, where X_i are chiral superfields of the model. This ensures that the D-term part of the potential vanishes[5] along the direction $\langle\chi_1\rangle = ... = \langle\chi_n\rangle = \varphi$ (χ_i are scalar components of X_i). This corresponds to a two-dimensional subspace represented by a complex field φ. A flat direction VEV spontaneously breaks gauge symmetries and gives (SUSY conserving) masses to the gauge bosons/gauginos similar to the Higgs mechanism in electroweak symmetry breaking [48, 61, 63, 68, 69]. The induced masses for gauge/gaugino fields are $\sim \alpha^{1/2}|\varphi|$ (we recall that α is a gauge fine structure constant). Similarly, a flat direction VEV induces (SUSY conserving) masses $\sim h|\varphi|$ for those fields that have superpotential couplings to φ (h is a Yukawa coupling). Therefore all fields that are coupled to a flat direction obtain very large masses.

The flat directions are massless if SUSY is exact, but they are lifted when SUSY is broken (which is assumed to happen at a scale of the order of TeV), as a result of which they get a mass $m_\varphi \sim \mathcal{O}(\mathrm{TeV})$. Provided that $m_\varphi \ll H_{\rm inf}$, $H_{\rm inf}$ being the Hubble expansion rate during inflation, the flat direction can acquire a large VEV by the virtue of quantum jumps during inflation, see the discussion in Refs. [65, 66]. This can dramatically alter the post-inflationary history of the universe as we will see in the next subsections [6].

[4] Note that the cubic term is required for a complete decay of the inflaton field.

[5] Since the total SM charge of a gauge-invariant monomial is zero by definition, the D-term potential involving only the fields used to build the monomial will also vanish since it is proportional to the sum of the charges.

[6] The development of large VEVs requires that the flat directions do not obtain positive Hubble-induced supergravity corrections during inflation. This problem can be avoided, for example, by considering non-minimal Kahler potentials.

C. Perturbative Decay

Consider a flat direction φ that has Yukawa couplings to the inflaton decay products χ. This happens, for example, for MSSM flat directions that are made of squark and/or slepton fields with χ being a MSSM Higgs field, see Eq. (56) (for details, see [61]). This results in the following term in the scalar potential:

$$V \supset h^2|\varphi|^2\chi^2, \tag{57}$$

where h denotes a Yukawa coupling. Note that the first generation of leptons and quarks have a Yukawa coupling $\sim \mathcal{O}(10^{-5})$, while the rest of the SM Yukawa couplings are $> 10^{-4}$. Since $|\varphi|$ is virtually frozen while $m_\varphi < H < H_{\rm inf}$ it is only when $H \simeq m_\varphi$ that the flat direction starts its oscillations. Since the field is complex, typically an elliptical trajectory with an $\mathcal{O}(1)$ eccentricity will result [63]. Hence, $|\varphi|$ will redshift as $|\varphi| \propto H^{-1}$.

While the flat direction has a large amplitude, the induced mass of the inflaton decay products obtained via (57) will lead to the inflaton decay being kinematically forbidden as long as $h|\varphi| \geq m/2$. There are thus two criteria for perturbative inflation decay. First of all, the decay of the inflaton into χ particles must be kinematically allowed which will become possible once the induced χ mass drops below the inflaton mass m. Taking into account the fact that once H falls below the value m_φ the field amplitude of φ decreases linearly in H we find that the kinematic decay becomes possible once $H < \left(\frac{m}{h\varphi_0}\right) m_\varphi$, where φ_0 is the initial VEV of the flat direction. A second condition for perturbative inflaton decay to occur is that $H < \Gamma$, Γ being the rate for perturbative decay of the inflaton. Thus, the inflaton cannot decay until the Hubble rate has decreased to a value $H_{\rm dec}$ given by:

$$H_{\rm dec} = \min \left[\left(\frac{m}{h\varphi_0}\right) m_\varphi, \quad \Gamma\right]. \tag{58}$$

If φ_0 is sufficiently large, then we can have $H_{\rm dec} \ll \Gamma$. This happens if $\varphi_0 > h^{-1}\frac{m_\varphi m}{\Gamma}$. Flat directions can therefore significantly delay inflaton decay on purely kinematical grounds.

D. Non-perturbative Decay

In order to understand the preheating dynamics in the presence of flat directions, we consider the governing potential that is obtained from Eqs. (56,57):

$$V = \frac{1}{2}m^2\phi^2 + g^2\phi^2\chi^2 + \frac{g}{\sqrt{2}}m\phi\chi^2 + h^2|\varphi|^2\chi^2. \tag{59}$$

As mentioned in the previous section, we generically have $h > 10^{-4}$, and g can be as large as $\sim \mathcal{O}(1)$. After mode decomposition of the field χ, the energy of the mode with