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AXINO MASS*

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ABSTRACT

The mass of the axino is computed in realistic supersymmetric extensions of the standard model. It is found to be strongly model dependent and can be as small as a few keV but also as large as the gravitino mass. Estimates of this mass can only be believed once a careful analysis of the scalar potential has been performed.

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While the standard model of electroweak interactions is in complete agreement with all experimental findings, still much effort has been devoted to construct various generalizations. Supersymmetric extensions have attracted much attention in the last decade.[1] Realistic models of this type consider local supersymmetry (supergravity), spontaneously broken in a so-called hidden sector at a mass scale of $M_S \approx 10^{11}$GeV. The induced SUSY-breaking mass scale in the observable sector is given by a value of order of the gravitino mass $m_{3/2} \sim M_S^2/M_P$, where M_P denotes the Planck mass. The breakdown scale of the weak interaction gauge symmetry $SU(2) \times U(1)$ is then closely related to the scale of SUSY-breakdown in the observable sector, thus explaining the smallness of M_W compared to M_P, once we understand the mechanism for the breakdown of supersymmetry.

In the minimal supersymmetric extension of the standard model, however, there is one dimensionful parameter which is not proportional to $m_{3/2}$, the Higgs mass term $\mu H\bar{H}$ in the superpotential. Such a term is allowed by supersymmetry and it remains to be understood why also μ should take values in the desired energy range. One suggestion to achieve this starts with the consideration of the axionic generalization of the standard model.[2] The decay constant F_a of the invisible axion should lie in the range given by M_S and it was argued that such a coincidence cannot be accidental. In fact, one can easily construct models, in which μ is generated dynamically to be in the range of the gravitino mass.[3]

Of course, the prime motivation to consider the axionic generalization of the standard model is the quest for a natural solution of the strong CP-problem. A supersymmetric realization of this mechanism is most easily achieved in a model with several Higgs supermultiplets.[4] Instead of just the pseudoscalar axion, such a model now possesses a full axion supermultiplet. This contains the *axino*, the fermionic partner of the axion, as well as the *saxino*, the scalar partner of the axino. Although these particles are very weakly interacting, they might nonetheless lead to important astrophysical and cosmological consequences. The stability of stars and the observed energy density of the universe, as is well known, restrict the decay constant of the invisible axion to a small window.[5] Furthermore, a very light saxino might give rise to new long range interactions which are incompatible with present observations. In a model of broken supersymmetry, however, one would usually expect the saxino to receive a mass of order of the gravitino in the 100 GeV to TeV range, avoiding such unpleasant phenomena as a fifth force.

The possible consequences of the presence of an axino have not been considered in detail until recently.[6] Given the weakness of its interactions, of course, we would rather expect to find only indirect manifestations of the existence of such a particle. In fact, up to now only the effects of the (possibly) stable axino on the total energy density of the universe have been studied. A stable axino in a certain mass range could lead to overcritical energy density, and the corresponding axion models are therefore ruled out. If, on the other hand, the axino has a mass of a few keV, it is itself an interesting candidate[6] for a source of dark matter. In fact, such a particle is up to now the only well motivated candidate for so-called *warm* dark matter. It still remains to be seen, however, whether warm dark matter can lead to a satisfactory cosmological model including questions about large scale structure formation.

In any case, the existence of a (light) axino might have important consequences in any supersymmetric extension of the standard model. Many of these models contain discrete symmetries (R-parity in the simplest case) that allow only pair production of the new supersymmetric particles. In these cases, there exists a lightest supersymmetric particle (called LSP) which is stable. In the minimal model one usually considers such weakly interacting massive particles (WIMPS, an example can be found in the photino) as a possible source of *cold* dark matter. In the presence of an axion supermultiplet, it could very well happen, that the axino is the LSP and thus render the WIMP unstable, at least on cosmological time scales.

While other properties of the axino seem not to be so important for our discussion, its mass is a crucial parameter and a careful analysis is required. We shall see in the following that this value is strongly model dependent. Before we discuss these questions in detail, let us remark, however, that, in general supergravity models, there are some *natural* values such a mass can have. One of them could be the mass of the gravitino that sets the scale of SUSY breakdown in the observable sector. But this is not the only possibility. The axino could very well be much lighter. In fact, models based on supergravity contain a very small dimensionless parameter

$$\eta = \frac{M_S}{M_P} \approx 10^{-8}. \tag{1}$$

The gravitino mass is then given by $m_{3/2} \sim \eta^2 M_P$ and the natural values for the mass of the axino at the tree level are given by

$$m_{\tilde{a}} \sim \eta^k M_P \tag{2}$$

with $k \geq 2$.

It is interesting to observe, that in the case of $k = 3$ this mass is in the region of 1 to 10 keV, leading to a critical mass density of the universe. In models of global supersymmetry one obtains similar estimates. Here $\eta \sim M_W/F_a$ where M_W denotes the scale of weak interaction breakdown.[6] These values coincide, since M_S and F_a are so close to each other.

The task of determining the mass of the axino in a given model now boils down to the question about the power k appearing in (2). For large k, of course, also radiative corrections to $m_{\tilde{a}}$ have to be taken into account.

Let us start our discussion in the framework of globally supersymmetric models. Although the construction of supersymmetric generalizations nowadays exclusively considers locally supersymmetric (supergravity) models, we can still learn a lot from the simpler models based on global SUSY. In the present example we can see quite easily, why it makes sense to consider the possibility of a very small axino mass. In the case of unbroken supersymmetry, the whole axion supermultiplet will remain degenerate at the mass given by the anomaly, which we shall neglect in the following. Thus the mass of the axino and the saxino have to be proportional to the scale of SUSY breakdown represented by the vacuum expectation value of an auxiliary field F_G (this is the auxiliary field of the goldstino multiplet)†. The mass splitting of the chiral supermultiplet is determined by the coupling of its members[1] to F_G. The axion is protected by a symmetry and does not receive a mass in the presence of SUSY breakdown. The scalar saxino couples in general to F_G and will thus obtain a mass of the order $g\,\langle F_G\rangle$, where g is the coupling to the goldstino multiplet*. This is the reason why one usually assumes the saxino to be heavy. In the case of the axino the situation is similar, but different. Again its mass is determined by the coupling to F_G, but the auxiliary field has canonical dimension two. A mass term for the axino $\tilde{a}\tilde{a}F_G$ is of dimension five and there are *no renormalizable contributions to the mass of the axino.* In a model with an (invisible) axion we have as additional dimensionful parameter the axion decay constant F_a of order of 10^{11} GeV and we therefore expect a small axino mass $m_{\tilde{a}} \sim \frac{F_G}{F_a}$ as was demonstrated in ref. 7.

It remains to be seen, how these results generalize once we consider models based on supergravity. The reason why one nowadays primarily considers these models is the fact that in models based on spontaneously broken global SUSY a universal mass shift for the scalar partners of quarks and leptons is not possible. We have mentioned that already in connection with the discussion of the mass of the saxino. This fact holds for a large class of models and can be succinctly summarized by the value of $STrM^2$, the supertrace of the square of the mass matrix. These results suggest that in realistic models the masses of the scalars, and thus also the mass of the saxino, are pushed up to a value beyond the reach of present experiments. In the case of the axino such a general statement cannot be made. The authors of ref. 6 assume (in order to avoid the murky depths of supergravity theory as they say), that the globally supersymmetric results carry

† F_G is in general a combination of the auxiliary fields of gauge and chiral supermultiplets.

$*$ Actually, in many models based on global supersymmetry this coupling can be very small and even vanish at tree level. These vanishing scalar masses, however, are the reason, why globally supersymmetric models do not lead to a realistic generalization of the standard model. We shall come back to this point later.

over to the supergravity case. We shall see in the following that, in general, such an assumption is not necessarily correct. A similar conclusion has been obtained by Goto and Yamaguchi.[8] Their result seems to to imply, however, that a small mass of the axino requires a special form of the kinetic terms. We analyze this issue in a more general way and see that, independent of the choice of the kinetic terms, small (and also large) axino masses are possible, dependent on other properties of the theory. We also investigate the question of the axino mass in those models that might be found as the low energy limit of string theory.

The scalar sector of a supergravity theory is completely specified by the Kahler potential $G(\Phi^j, \Phi^*_j)$ where Φ collectively denotes the chiral superfields. The scalar kinetic terms are given by the second derivative $G^i_j = \partial^2 G/\partial\Phi^j \partial\Phi^*_i$ and one often splits $G(\Phi, \Phi^*) = K(\Phi, \Phi^*) + \log|W(\Phi)|^2$ where the superpotential $W(\Phi)$ is a holomorphic function of Φ. The scalar potential is given by[1]

$$V = -\exp G\left[3 - G_i(G^{-1})^i_j G^j\right]. \tag{3}$$

We are interested in the mass spectrum of the theory once supersymmetry is broken spontaneously, which leads to a nontrivial value of the gravitino mass $m^2_{3/2} = \exp(G)$ †. Masses of the scalar particles can then be read off from the second derivative of the potential at the minimum. For the fermions we obtain

$$M_{ij} = \exp(G/2)\left[G_{ij} + \frac{1}{3}G_i G_j - G_k(G^{-1})^k_l G^l_{ij}\right], \tag{4}$$

where we have removed the contribution to the mass of the gravitino. We also have to respect the constraint from the anomalous $U(1)$-symmetry

$$\sum_i \left[q_i \Phi^i G_i - q_i \Phi^*_i G^i\right] = 0, \tag{5}$$

where q_i is the $U(1)_{PQ}$-charge of Φ_i.

† We assume vanishing vacuum energy, thus $\langle V\rangle = \langle V_i\rangle = 0$ at the minimum.

Fields in the observable (hidden) sector shall be denoted by y_i (z_i), respectively, and we shall assume the superpotential to split: $W(\Phi_i) = h(z_i) + g(y_i)$. In our examples we use the well known case $h(z) = m^2(z+\beta)$ for simplicity. Let us start our discussion with a special choice $G_i^j = \delta_i^j$, usually referred to as minimal kinetic terms. The scalar potential then reads

$$V = \exp(K/M^2)\left[|h_z + \frac{z^*W}{M^2}|^2 + |g_i + \frac{y_i^*W}{M^2}|^2 - \frac{3|W|^2}{M^2}\right]. \tag{6}$$

Within this framework Goto and Yamaguchi[8] have argued that the axino mass is as large as the gravitino mass. Let us see how this works using their superpotential $g_1 = \lambda(AB - f^2)Y$, where f is a constant and A, B and Y are fields. Minimizing (6) we find the following (approximate) vacuum expectation values (vevs): $A = B \approx f$ and $Y \approx m^2/M$, while the z vev remains undisturbed*. Actual values for f and m should lie in the range of 10^{11} GeV. Denoting the fermions in the chiral supermultiplets by χ^i, the axino is found to be the linear combination $\chi^a = (\chi^A - \chi^B)/\sqrt{2}$ and it receives a mass of order of the gravitino mass $m_{\tilde{a}} \sim m_{3/2} \sim m^2/M$.

Is this now a generic property of models with minimal kinetic terms? We shall see that the answer is no by inspecting a second example with superpotential $g_2 = \lambda(AB - X^2)Y + \frac{\lambda'}{3}(X-f)^3$ including a new singlet chiral superfield X. Minimization of the potential now becomes more complicated since the condition $\langle V\rangle = 0$ leads to a shift in the vev of the hidden sector fields. The easiest way to discuss the potential is by expanding it in powers of m/M. To lowest order one obtains the globally supersymmetric result for the observable sector. In each order one then has to adjust the vacuum energy to zero, and in the present example the inclusion of the terms of order m^2/M^2 require a shift of $\langle z\rangle$. In the previous example it was sufficient to just consider the expansion up to first order. One still obtains vevs of A, B similar to those of the previous example, but the presence of the field X has important consequences on the axino mass; in fact here one obtains $m_{\tilde{a}} \sim m^3/M^2$.

* Observe that in the case of global supersymmetry the minimum is found at $\langle Y\rangle = 0$.

This example shows that the mass of the axino depends strongly on the model and the special form of the superpotential. It also shows that in models with minimal kinetic terms the mass of the axino not necessarily needs to be as large as the gravitino mass, contrary to the impression given in ref. 8. In particular, masses of the axino in the range of a few keV can be obtained also in this framework.

Let us next consider those supergravity models that have a structure similar to those that appear in the low-energy limit of string theories. The Kahler potential is given by[9]

$$K = -\log(S + S^*) - 3\log(T + T^* - C^i C_i^*), \tag{7}$$

where S denotes the dilaton superfield, T represents the moduli and C_i the matter fields. We shall assume the superpotential of the form $W = W(S) + W(C_i)$, postponing a discussion of the implications of moduli dependence. The term $W(S)$ is assumed to appear as a result of gaugino condensation in the underlying string model, and is crucial for the process of supersymmetry breakdown. For a review and details see ref. 10. The scalar potential of the theory defined in this way

$$V = \exp(G)\left[|G_S|^2(S + S^*) + |W_i|^2\right] \tag{8}$$

is positive with a minimum at $\langle V \rangle = 0$; the dilaton adjusts its vev to cancel any possible contribution to the vacuum energy. Supersymmetry is broken spontaneously through a nontrivial vev of the auxiliary field of the dilaton supermultiplet $F_T \sim \exp(G)G_T$, while $F_S = 0$. The only problem with the potential is the fact that the vevs of the moduli are not determined and thus the vacuum is highly degenerate. Let us nonetheless discuss this simplified example first. The minimum of (8) is found at $\langle G_S \rangle = \langle W_i \rangle = 0$, where $W_i = \partial W / \partial C_i = 0$ coincides with the solution obtained in the case of global supersymmetry, independent of the special form of the superpotential. In our case we require nontrivial vevs $\langle C_i \rangle = v_i$ for at least one of the charged scalar fields. The axino is then given by $\tilde{a} = \sum_i (\frac{q_i v_i}{v})\chi^i$, where $v = \sqrt{\sum_i q_i^2 v_i^2}$ should take a value of order of m in the 10^{11} GeV range. The goldstone fermion is given by $\eta \sim G_T \chi^T + G_i \chi^i$, with $G_T = -3/\Delta$, $G_i = 3v_i/\Delta$ and $\Delta = T + T^* - C^i C_i^*$. One thus obtains $\sum q_i v_i G_i = 0$ for the axino to be orthogonal to the goldstino. Fermion masses can now be computed according to (4) in a straightforward manner. This gives e.g. $M_{TT} \sim G_{TT} + \frac{1}{3} G_T G_T + \frac{2}{\Delta} G_T = 0$, since $G_{TT} = 3/\Delta^2$ and $G_T = -3/\Delta$. Also the terms mixing T- and i-components vanish as well as $M_{aj} = \sum_i (\frac{q_i v_i}{v}) M_{ij}$ because of the constraint (5). Thus all these fermions including the axino remain *massless*.

One could have expected such a result from the outset because of the fact that models with kinetic terms of the structure (7) are very closely related to globally supersymmetric models. We have confirmed that above finding $W_i = 0$ at the minimum, the globally supersymmetric solution. Thus one might obtain a light axino in a natural way.[8] But this is probably not the whole story. The other fermions remain massless as well and, more importantly, also the scalar particles like the saxino remain massless at tree level. At the present stage of the discussion we can conclude that this model not only shares the desirable features of globally supersymmetric models but also the more problematic ones. Observe that in the models based on minimal kinetic terms the scalar particles and thus also the saxino received a large mass of order of $m_{3/2}$.

Again the question arises whether in models with Kahler potential as in (7) one always obtains a light axino. Unfortunately this question cannot yet be answered definitely. One way to proceed is to compute radiative corrections and see how axino and saxino masses are shifted.[8] We would like to argue, however, that this is not necessarily the correct way to attack this problem. After all the potential given in (8) has still a large vacuum degeneracy and many massless scalars and thus is unstable under small changes of the parameters. In fact, naively including radiative corrections might destabilize the potential in such a way that it becomes unbounded from below. As long as we do not know the correct position of the minimum we can not really be sure that our estimate of the axino mass is reliable. This can be demonstrated quite easily in the framework of explicit models. We have seen this in our discussion of the models with minimal kinetic terms comparing those with superpotentials g_1 and g_2. Although there the potential is less unstable the actual value of the axino mass strongly depends on details of the potential. Similar things will happen also in models with nonminimal kinetic terms.

In addition we know that the potential as given in (8) is incomplete. In a first step one should include the moduli-dependent contributions to the superpotential. Unfortunately the incorporation of such a dependence in W leads to enormous complications. The potential is no longer positive definite and nobody succeeded yet to find a satisfactory minimum with broken supersymmetry and a vanishing cosmological constant. As long as such a result is missing, any reliable computation of the axino mass in such models is impossible. Unfortunately this is also true in those models with a composite axino that constituted our prime motivation to study these questions in detail.[11,12] This does not mean that the axino cannot be light. In fact our discussion of the models with minimal kinetic terms has demonstrated that light axinos could actually exist. We want to stress here that any estimate of axino masses is unreliable as long as a detailed calculation of the underlying potential has not been performed.

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REFERENCES

1. H. P. Nilles, *Physics Reports* **110** (1984), 1
2. Jihn E. Kim, *Physics Reports* **150** (1987), 1
3. J. E. Kim and H. P. Nilles, *Phys. Lett.* **B138** (1984), 150
4. H. P. Nilles and S. Raby, *Nucl. Phys.* **B198** (1982), 102
5. G. G. Raffelt, *Physics Reports* **198** (1990), 1
6. K. Rajagopal, M. S. Turner and F. Wilczek, *Nucl. Phys.* **B358** (1991), 447
7. K. Tamvakis and D. Wyler, *Phys. Lett.* **B112** (1982), 451
8. T. Goto and M. Yamaguchi, Tohoku University preprint TU-378 (1991)
9. E. Witten, *Phys. Lett.* **B155** (1985), 151
10. H. P. Nilles, *Int. Journ. of Mod. Phys.* **A5** (1990), 4199
11. J. E. Kim and H. P. Nilles, *Phys. Lett.* **B263** (1991), 79
12. E. J. Chun, J. E. Kim and H. P. Nilles, *Nucl. Phys.* **B370** (1992), 105